Scherk-Schwarz Supersymmetry Breaking for Quasi-localized Matter Fields and Supersymmetry Flavor Violation

Hiroyuki Abe*, Kiwoon Choi†, Kwang-Sik Jeong‡, Ken-ichi Okumura§
Department of Physics,
Korea Advanced Institute of Science and Technology
Daejeon 305-701, Korea
(Dated: February 7, 2008)

We examine the soft supersymmetry breaking parameters induced by the Scherk-Schwarz (SS) boundary condition in 5-dimensional orbifold field theory in which the quark and lepton zero modes are quasi-localized at the orbifold fixed points to generate the hierarchical Yukawa couplings. In such theories, the radion corresponds to a flavon to generate the flavor hierarchy and at the same time plays the role of the messenger of supersymmetry breaking. As a consequence, the resulting soft scalar masses and trilinear A-parameters of matter zero modes at the compactification scale are highly flavor-dependent, thereby can lead to dangerous flavor violations at low energy scales. We analyze in detail the low energy flavor violations in SS-dominated supersymmetry breaking scenario under the assumption that the compactification scale is close to the grand unification scale and the 4-dimensional effective theory below the compactification scale is given by the minimal supersymmetric standard model. Our analysis can be applied to any supersymmetry breaking mechanism giving a sizable F-component of the radion superfield, e.g. the hidden gaugino condensation model.

I. INTRODUCTION

Supersymmetry (SUSY) is one of the prime candidates for new physics beyond the standard model (SM) [1]. An important issue in supersymmetric theories is to understand how SUSY is broken in low energy world. It has been known that theories with compact extra dimension provide an attractive way to break SUSY by imposing nontrivial boundary conditions on the field variables. This mechanism which has been proposed originally by Scherk and Schwarz (SS) [2] can be interpreted as a SUSY breaking induced by the auxiliary component of higher dimensional supergravity (SUGRA) multiplet [3]. Extra dimension can provide also an attractive mechanism to generate hierarchical Yukawa couplings [4]. The quark and lepton fields can be quasi-localized in extra dimension, and then their 4-dimensional (4D) Yukawa couplings involve the wavefunction overlap factor $e^{-M\pi R}$ where M is a combination of mass parameters in higher dimensional theory and R is the size of extra dimension. This would result in hierarchically different Yukawa couplings even when the fundamental mass parameters have the same order of magnitudes.

A simple and natural theoretical framework for the quasi-localization of matter zero modes is supersymmetric 5D orbifold field theory. If a matter hypermultiplet in 5D orbifold SUGRA has a non-zero gauge charge for the graviphoton and/or for an ordinary $U(1)_{FI}$ vector multiplet which has non-zero boundary Fayet-Iliopoulos (FI) terms [5], it obtains a non-zero 5D kink mass $M\epsilon(y)$ where $\epsilon(y) = \pm 1$ is the periodic sign function on S^1/Z_2 whose fundamental domain is given by $0 \leq y \leq \pi$. In the presence of such kink mass, the matter zero mode becomes quasi-localized at one of the orbifold fixed points $y = 0, \pi$ with a wavefunction given by $e^{-MR|y|}$. A 5D orbifold SUGRA provides also a simple theoretical framework for the SS SUSY breaking. The theory admits a continuous twist of $SU(2)_R$ boundary condition under the discrete shift $y \to y + 2\pi$ which would break the $N = 1$ SUSY survived from the Z_2-orbifolding [3].

In this paper, we wish to examine some physical consequences of implementing the quasi-localization of matter zero modes and the SS SUSY breaking simultaneously within 5D orbifold field theories, particularly the flavor structure of soft parameters and the resulting low energy flavor violations. In section 2, we first discuss some features of 5D orbifold SUGRA related to the quasi-localization of matter fields and also the SS SUSY breaking. We then compute the soft parameters of quasi-localized matter fields induced by the SS boundary condition in generic 5D orbifold SUGRA. We show explicitly that the zero mode soft parameters from the SS boundary condition are same as the ones induced by the radion F-component in 4D effective SUGRA, and thus our analysis applies to any SUSY breaking mechanism giving a sizable F-component of the radion superfield [6], e.g. the hidden gaugino condensation model. This means that the radion superfield which corresponds to a flavon for the Yukawa hierarchy plays the role of the

* abe@hep.kaist.ac.kr
† kchoi@hep.kaist.ac.kr
‡ ksjeong@hep.kaist.ac.kr
§ okumura@hep.kaist.ac.kr

messenger of SUSY breaking. As a consequence, the resulting soft scalar masses and trilinear A parameters of matter zero modes at the compactification scale are highly flavor-dependent, thereby can lead to dangerous flavor violation at low energy scales. In particular, the predicted shape of soft parameters at the compactification scale indicates that the compactification scale should be much higher than the weak scale in order for the model to be phenomenologically viable.

In 5D orbifold SUGRA, 5D kink masses $M_I\epsilon(y)$ responsible for quasi-localization have quantized-values if the graviphoton and/or $U(1)_{FI}$ gauge charges are quantized. An important feature of the SS SUSY breaking is that, if the kink masses are quantized, the resulting soft scalar masses and the trilinear scalar couplngs (divided by the corresponding Yukawa couplings) at the compactification scale are quantized also in the leading approximation. This feature provides a natural mechanism to suppress dangerous flavor violations since the flavor violating amplitudes appear in a form $f(M_I) - f(M_J)$, thus are canceled when some of the quantized kink masses are degenerate.

In section 3, we analyze in detail the resulting low energy flavor violations under the assumption that the compactification scale M_c is close to the grand unification scale $M_{GUT} \sim 2\times 10^{16}$ GeV and the 4D effective theory below M_c is given by the minimal supersymmetric standard model (MSSM). We find that many of the low energy flavor violations are appropriately suppressed, however generically ϵ_K and $\mu \to e\gamma$ can be too large if the SS boundary condition is the major source of SUSY breaking. As summarized in Table I, if either the $SU(2)_L$-doublet lepton kink masses or the $SU(2)_L$-singlet lepton kink masses are flavor-independent, the $\mu \to e\gamma$ bound can be satisfied for a reasonable range of the involved continuous parameters. Tables II−III contain the predictions for lepton flavor violating processes of the models which can satisfy the $\mu \to e\gamma$ constraint with a mild tuning of the involved parameters. Different choices of the lepton kink masses predict different patterns of lepton flavor violations. In particular, when the $SU(2)_L$-doublet lepton kink masses are degenerate, the predicted chirality structure of decay modes is opposite to the other case with degenerate $SU(2)_L$-singlet lepton kink masses which has the same chirality structure as the lepton flavor violating decays in seesaw models [7, 8, 9]. For ϵ_K, it is more difficult to make the SUSY contribution small enough since the model is constrained to yield the correct CKM mixing angles as well as the correct quark mass hierarchy. Again a possible option is that the kink masses of the $SU(2)_L$ singlet down quarks are flavor-independent. However in this case, in order to produce the correct quark mass eigenvalues and CKM mixing angles, one needs to assume that some boundary Yukawa couplings are abnormally large (or small) by a factor of $4 \sim 5$ ($0.2 \sim 0.3$) compared to the values suggested by the naive dimensional analysis. Table IV summarizes the SUSY contributions to ϵ_K for some choices of the quark kink masses. We finally discuss the SUSY contributions to other flavor-violating amplitudes in SS SUSY breaking scenario, e.g. the $b \to s\gamma$ rate, ϵ'/ϵ_K and the K^0-$\bar{K}^0$ and B^0-$\bar{B}^0$ mass differences, which turn out to be either well below or at most comparable to the SM contributions.

II. SUPERSYMMETRY BREAKING BY BOUNDARY CONDITION FOR QUASI-LOCALIZED MATTER FIELDS IN 5D ORBIFOLD SUPERGRAVITY

In this section, we first discuss some features of 5D orbifold SUGRA related to the quasi-localization of matter fields and also the SS SUSY breaking by boundary condition. We then compute the soft parameters of quasi-localized matter fields induced by the SS boundary condition and compare the results with the radion-mediated soft parameters in 4D effective SUGRA.

Let us consider a generic SUGRA-coupled 5D gauge theory on S^1/Z_2. The action of the theory is given by [10]

$$\begin{aligned} S \;=\; \int d^5x\sqrt{-G}\,\Bigg[\frac{1}{2}\left(\mathcal{R} + \bar{\Psi}^i_M\gamma^{MNP}D_N\Psi_{iP} - \frac{3}{2}C_{MN}C^{MN} - \frac{3}{2}k\epsilon(y)\bar{\Psi}^i_M\gamma^{MN}\Psi_{iN} - 12k^2 + ...\right) \\ +\; \frac{1}{g^2_{5a}}\left(-\frac{1}{4}F^{aMN}F^a_{MN} + \frac{1}{2}D_M\phi^aD^M\phi^a + \frac{i}{2}\bar{\lambda}^{ai}\gamma^MD_M\lambda^a_i + \frac{1}{2}k\epsilon(y)\bar{\lambda}^{ia}\lambda^a_i - 4k^2\phi^a\phi^a + ...\right) \\ +\; \left(|D_Mh^i_I|^2 + i\bar{\Psi}_I\gamma^MD_M\Psi_I + iM_I\epsilon(y)\bar{\Psi}_I\Psi_I + (M_I^2 \pm kM_I - \frac{15}{4}k^2)|h^i_I|^2 + ...\right)\Bigg]\,, \end{aligned} \tag{1}$$

where $\mathcal{R}$ is the Ricci scalar of the 5D metric G_{MN}, Ψ^i_M ($i = 1,2$) are $SU(2)_R$-doublet 5D gravitino, $C_{MN} = \partial_MB_N - \partial_NB_M$ is the graviphoton field strength. The fields $(\phi^a, A^a_M, \lambda^{ia})$ are 5D real scalar, vector and $SU(2)_R$-doublet gaugino constituting a vector multiplet, (h^i_I, Ψ_I) are $SU(2)_R$-doublet complex hyperscalar and Dirac hyperino constituting a hypermultiplet, and $\epsilon(y) = \pm 1$ is the periodic sign function on S^1/Z_2 satisfying $\epsilon(y) = \epsilon(y + 2\pi) = -\epsilon(-y)$. Here we set the 5D Planck mass $M_5 = 1$, and the ellipses include appropriate boundary actions.

In 5D orbifold SUGRA, the bulk and boundary cosmological constants and also the hyperino kink masses appear in connection with the graviphoton gauge couplings [11]:

$$D_Mh^i_I = \nabla_Mh^i_I - i\Big(\frac{3}{2}k(\sigma_3)^i_j - c_I\delta^i_j\Big)\epsilon(y)B_Mh^j_I,$$

$$D_M\Psi_I = \nabla_M\Psi_I + ic_I\epsilon(y)B_M\Psi_I,$$
$$D_M\lambda^{ia} = \nabla_M\lambda^{ia} - i\frac{3}{2}k(\sigma_3)^i_j\epsilon(y)B_M\lambda^{aj}, \quad (2)$$

where ∇_M contains other gauge couplings. For instance, the Z_2-odd $U(1)_R$ gauge coupling of the graviphoton, $\frac{3}{2}k\epsilon(y)\sigma_3$, is associated with the bulk and boundary cosmological constants:

$$-6k^2 + 6k\frac{(\delta(y) - \delta(y-\pi))}{\sqrt{G_{55}}}$$

which leads to the warped Randall-Sundrum geometry. As for the hyperino kink mass $M_I\epsilon(y)$, another possible origin is a $U(1)_{FI}$ vector multiplet whose scalar component develops a kink-type vacuum expectation value due to the boundary FI terms [5]. Including this FI contribution, the (effective) hyperino kink mass is given by

$$M_I = c_I + q_I\xi_{FI}\,, \quad (3)$$

where q_I is the $U(1)_{FI}$ charge of Ψ_I and ξ_{FI} is the FI coefficient. Throughout this paper, we will assume that the $U(1)$ gauge charges c_I, q_I are *quantized*, and thus the hyperino kink masses are quantized also.

In 5 dimension, orbifolding the theory corresponds to imposing the boundary condition

$$\Phi(-y) = Z\Phi(y)\,, \quad \Phi(y+2\pi) = \Omega\Phi(y)\,, \quad (4)$$

for generic 5D field Φ, where Z and Ω satisfy the consistency conditions

$$Z^2 = 1\,, \quad Z\Omega = \Omega^{-1}Z. \quad (5)$$

Equivalently, one can impose the parity boundary condition at each fixed point:

$$\Phi(-y) = Z\Phi(y)\,, \quad \Phi(-y') = Z'\Phi(y') \quad (Z^2 = Z'^2 = 1)\,, \quad (6)$$

where $y' = y - \pi$ and $Z' = Z\Omega$. In case that Z and Ω commute to each other, the consistency condition (5) implies $\Omega^2 = 1$, and thus Z and Ω can be simultaneously diagonalized to have eigenvalues $Z = \pm 1$ and $\Omega = \pm 1$. On the other hand, for the case of SS boundary condition, Z and Ω do not commute, so Ω can have a continuous value. Still one can adopt a field basis for which Z is diagonal, while Ω (and thus $Z' = Z\Omega$) is not diagonal in general.

The orbifolding boundary condition (4) should be consistent with all gauge symmetries of the theory, including the Z_2-odd graviphoton gauge transformation:

$$B_M \;\rightarrow\; B_M + \partial_M\Lambda, \quad \Phi \;\rightarrow\; e^{i\epsilon(y)Q\Lambda}\Phi, \quad (7)$$

where Q is a constant charge matrix and the transformation function satisfy $\Lambda(y) = \Lambda(y+2\pi) = -\Lambda(-y)$ in order to be consistent with 5D local SUSY. In order for $D_M\Phi = (\nabla_M - i\epsilon(y)QB_M)\Phi$ to have a consistent boundary condition, both Z and Ω should commute with Q:

$$ZQ = QZ, \quad \Omega Q = Q\Omega \quad (8)$$

which correspond to additional consistency condition for orbifolding boundary conditions.

5D orbifold SUGRA admits a continuous twist of $SU(2)_R$ boundary condition under the discrete shift $y \rightarrow y + 2\pi$, breaking the $N = 1$ SUSY survived from the Z_2-orbifolding. Let Z_R, Ω_R and Q_R denote the $SU(2)_R$ representations of Z, Ω and Q, respectively. To obtain 4D chiral fermion, one can always choose $Z_R = \sigma_3$, and then the graviphoton $U(1)_R$ charge is given by $Q_R = \frac{3}{2}k\sigma_3$ as in (2). The consistency condition (5) implies that a continuous SS twist can be written as $\Omega_R = \exp(i\vec{\omega}\cdot\vec{\sigma})$ where $\vec{\omega} = (\omega_1, \omega_2, 0)$. However, if $Q_R = \frac{3}{2}k\sigma_3$ is non-vanishing, there doesn't exist any non-trivial SS twist allowed by the consistency condition (8). In other words, a continuous SS SUSY breaking is not allowed in 5D orbifold SUGRA yielding a warped Randall-Sundrum geometry as has been noticed in [12]. Thus in the following, we will focus on the case that the graviphoton $U(1)_R$ charge vanishes:

$$Q_R = \frac{3}{2}k\sigma_3 = 0,$$

which gives a flat spacetime geometry

$$ds^2 = \eta_{\mu\nu}dx^\mu dx^\nu - R^2dy^2$$

and allows a continuous SS twist

$$\Omega_R = \exp(2\pi i\omega\sigma_2)$$

in an appropriate $SU(2)_R$ basis. Obviously, only the $SU(2)_R$-doublet gravitino, gauginos and hyperscalars are affected by this SS twist, e.g.

$$\begin{aligned}\lambda^{ai}(y+2\pi) &= \left(e^{2\pi i\omega\sigma_2}\right)^i_j \lambda^{aj}(y)\,,\\ h^i_I(y+2\pi) &= \left(e^{2\pi i\omega\sigma_2}\right)^i_j h^j_I(y)\,.\end{aligned} \tag{9}$$

It is convenient to write the 5D action (1) in $N=1$ superspace [13]. For the 5D SUGRA multiplet, we keep only the radion superfield

$$T = R + iB_5 + \theta\Psi_{5R} + \theta^2 F^T\,,$$

where $R=\sqrt{G_{55}}$ denotes the radius of the compactified 5-th dimension, and $\Psi_{5R} = \frac{1}{2}(1+\gamma_5)\Psi^{i=2}_{M=5}$. The relevant piece of the 5D action in $N=1$ superspace is given by

$$\int d^5x \left[\int d^4\theta\,\frac{T+T^*}{2}\left(H_I H_I^* + H_I^c H_I^{c*}\right) + \left(\int d^2\theta\, H_I^c\left(\partial_y + M_I T\epsilon(y)\right)H_I + \frac{1}{4g_{5a}^2} T W^{a\alpha} W^a_\alpha + \text{h.c.}\right)\right], \tag{10}$$

where W^a_α is the chiral spinor superfield for the 5D vector superfield

$$V^a = -\bar\theta\sigma^\mu\theta A^a_\mu - i\bar\theta^2\theta\lambda^a + i\theta^2\bar\theta\bar\lambda^a + \frac{1}{2}\theta^2\bar\theta^2 D^a\,,$$

and

$$\begin{aligned} H_I &= h^1_I + \theta\psi_I + \theta^2 F^1_I\,,\\ H^c_I &= h^{2*}_I + \theta\psi^c_I + \theta^2 F^{2*}_I\,,\end{aligned} \tag{11}$$

for $\lambda^a = \frac{1}{2}(1-\gamma_5)\lambda^{a1}$, $\psi_I = \frac{1}{2}(1-\gamma_5)\Psi_I$, and $\bar\psi^c{}_I = \frac{1}{2}(1+\gamma_5)\Psi_I$.

Here we are interested in 5D vector multiplets giving massless 4D gauge bosons, and also 5D hypermultiplets giving massless chiral 4D fermions. We thus consider the $N=1$ superfields satisfying the Z_2-boundary condition:

$$V^a(-y) = V^a(y), \quad H_I(-y) = H_I(y), \quad H^c_I(-y) = -H^c_I(y), \tag{12}$$

The zero mode equation for $\psi_I = \chi(x)\tilde\phi_{0I}(y)$ is given by

$$\left(\partial_y + M_I T\epsilon(y)\right)\tilde\phi_{0I} = 0,$$

yielding the zero mode wavefunction

$$\tilde\phi_{0I} \propto e^{-M_I T|y|},$$

which shows that the zero mode is quasi-localized at $y=0$ if $M_I>0$, and at $y=\pi$ if $M_I<0$.

In 5D orbifold SUGRA, Yukawa couplings can be introduced only through the boundary actions. For the orbifolding given by (6), the boundary action at $y=0$ is required to be invariant under the Z-even supercharges $\mathcal{Q}_Z$, while the boundary action at $y=\pi$ is invariant under the Z'-even supercharges $\mathcal{Q}_{Z'}$. Note that $\mathcal{Q}_Z$ and $\mathcal{Q}_{Z'}$ are related to each other by the SS twist: $\mathcal{Q}_Z = \Omega_R^{1/2}\mathcal{Q}_{Z'}$. Then in the presence of nontrivial SS twist, the boundary actions for Yukawa couplings can be written as

$$\int d^5x \int d^2\theta\left(\delta(y)\frac{1}{6}\lambda_{IJK} H_I H_J H_K + \delta(y-\pi)\frac{1}{6}\lambda'_{IJK} H'_I H'_J H'_K\right) + \text{h.c}\,, \tag{13}$$

where H_I and H^c_I are Z-even and Z-odd superfields, respectively, defined as (11) for the $N=1$ SUSY generated by $\mathcal{Q}_Z$, and H'_I and H'^c_I are Z'-even and Z'-odd superfields, respectively, for the $N=1$ SUSY generated by $\mathcal{Q}_{Z'}$. More explicitly,

$$\begin{aligned} H'_I &= h'^1_I + \theta\psi_I + \theta^2 F'^1_I,\\ H'^c_I &= h'^{2*}_I + \theta\psi^c_I + \theta^2 F'^{2*}_I,\end{aligned}$$

where

$$\begin{aligned} h_I'^1 &= \cos(\omega\pi) h_I^1 - \sin(\omega\pi) h_I^2, \\ h_I^{c\prime} &= \cos(\omega\pi) h_I^{2*} + \sin(\omega\pi) h_I^{1*}, \end{aligned}$$

and $F_I'^i$ can be determined by their equations of motion. Note that Z-odd H_I^c and Z'-odd $H_I'^c$ vanish at $y = 0$ and $y = \pi$, respectively, thus the boundary Yukawa operators involving H_I^c ($H_I'^c$) at $y = 0$ ($y = \pi$) vanish also.

Let us now compute the 4D Yukawa couplings, scalar masses and trilinear A parameters for the hypermultiplets obeying the boundary conditions (9) and (12). To this end, we analyze the Kaluza-Klein (KK) mass spectrum and wavefunctions of hypermultiplets, which has been done in [14]. The equation of motion for hyperscalar fields leads to the following KK wave equation:

$$\left(\frac{1}{R^2}\partial_y^2 + m_I^2 - M_I^2 + \frac{M_I}{R}\sigma_3 \partial_y \epsilon(y)\right) \tilde{\phi}_I(y) = 0, \tag{14}$$

where $\partial_y \epsilon(y) = 2(\delta(y) - \delta(y - \pi))$, and the KK wavefunction $\tilde{\phi}_I = (\tilde{\phi}_I^1, \tilde{\phi}_I^2)$ is defined as

$$h_I^i(x, y) = \phi_I(x) \tilde{\phi}_I^i(y)$$

for the 4D field $\phi_I(x)$ satisfying the on-shell condition

$$\partial_\mu \partial^\mu \phi_I(x) = m_I^2 \phi_I(x).$$

According to the boundary conditions (9) and (12), $\tilde{\phi}_I$ obeys

$$\tilde{\phi}_I^i(-y) = \left(\sigma_3\right)^i_{\ j} \tilde{\phi}_I^j(y), \quad \tilde{\phi}_I^i(y + 2\pi) = \left(e^{i2\pi\omega\sigma_2}\right)^i_{\ j} \tilde{\phi}_I^j(y). \tag{15}$$

It is then straightforward to find [14]

$$\begin{aligned} \tilde{\phi}_I^1 &= C_I \left(\cos \Delta_I y - \frac{M_I R}{\Delta_I} \sin \Delta_I y\right), \\ \tilde{\phi}_I^2 &= C_I \tan \omega\pi \left(-\cot \Delta_I \pi + \frac{M_I R}{\Delta_I}\right) \sin \Delta_I y, \end{aligned} \tag{16}$$

where Δ_I satisfies

$$\left(\frac{R M_I}{\Delta_I}\right)^2 + 1 = \frac{\sin^2 \omega\pi}{\sin^2 \Delta_I \pi}. \tag{17}$$

The corresponding KK mass eigenvalue is given by

$$m_I^2 = M_I^2 + \left(\frac{\Delta_I}{R}\right)^2, \tag{18}$$

and C_I is the normalization constant which can be determined by

$$2R \int_0^\pi dy \, \left(|\tilde{\phi}_I^1|^2 + |\tilde{\phi}_I^2|^2 \right) = 1\,.$$

Note that (16) represents the KK wavefunction over the fundamental domain $0 \le y \le \pi$. The KK wavefunction outside the fundamental domain can be determined by the boundary condition (15).

Once the KK wavefunctions are determined, the 4D couplings of the corresponding KK modes can be easily obtained. Let y_{IJK} and A_{IJK} denote the 4D Yukawa couplings and trilinear scalar couplings, respectively, of the *canonically normalized* 4D scalars ϕ_I and 4D fermions χ_I:

$$\int d^4x \left(\frac{1}{2} y_{IJK} \phi_I \chi_J \chi_K - \frac{1}{6} A_{IJK} \phi_I \phi_J \phi_K + \text{h.c.}\right). \tag{19}$$

Since the hyperinos are not affected by the SS twist, the hyperino KK wavefunctions correspond to the hyperscalar KK wavefunctions with $\omega = 0$, i.e.

$$h_I^i(x, y) = \phi_I(x) \tilde{\phi}_I^i(y), \quad \psi_I(x, y) = \chi_I(x) \tilde{\phi}_{0I}^i(y),$$

where $\tilde{\phi}^i_{0I} \equiv \tilde{\phi}^i_I|_{\omega=0}$. In our case, the 4D Yukawa couplings of KK modes are easily found to be

$$y_{IJK} \;=\; \int dy \left(\lambda_{IJK}\delta(y)\tilde{\phi}^1_I\tilde{\phi}^1_{0J}\tilde{\phi}^1_{0K} + \lambda'_{IJK}\delta(y-\pi)\tilde{\phi}'^1_I\tilde{\phi}^1_{0J}\tilde{\phi}^1_{0K} \right), \tag{20}$$

where

$$\tilde{\phi}'^i_I = \left(e^{i\pi\omega\sigma_2}\right)^i_j \tilde{\phi}^j_I \,.$$

Using the equation of motion for the auxiliary components F^i_I:

$$\begin{aligned} F^i_I &= \frac{1}{R}\left(i\sigma_2\right)^i_j \partial_y h^j_I - M_I\epsilon(y)\left(\sigma_1\right)^i_j h^j_I \,, \\ F'^i_I &= \frac{1}{R}\left(i\sigma_2\right)^i_j \partial_y h'^j_I - M_I\epsilon(y)\left(\sigma_1\right)^i_j h'^j_I \,, \end{aligned}$$

one can obtain also the trilinear soft scalar couplings as

$$\begin{aligned} A_{IJK} \;=\; & -\frac{1}{R}\int dy \left(\lambda_{IJK}\delta(y)\tilde{\phi}^1_I\tilde{\phi}^1_J\partial_y\tilde{\phi}^2_K + \lambda'_{IJK}\delta(y-\pi)\tilde{\phi}'^1_I\tilde{\phi}'^1_J\partial_y\tilde{\phi}'^2_K \right) \\ & + \Big(I \leftrightarrow K \Big) + \Big(J \leftrightarrow K \Big). \end{aligned} \tag{21}$$

The scalar masses and trilinear couplings discussed above can be interpreted as the parameters renormalized at the compactification scale $M_c \sim 1/R$. It is rather clear that the resulting soft parameters of quasi-localized matter zero modes can not be phenomenologically viable unless M_c is far above the weak scale M_W. Since the soft parameters have values of $\mathcal{O}(\omega M_c)$, one needs $|\omega| \ll 1$ to get the weak scale soft parameters for $M_c \gg M_W$. With this observation, in the following, we limit the discussion to the case of *small* SS parameter,

$$|\omega| \;\ll\; 1\,,$$

and compute the soft parameters of the canonically normalized hyperscalar zero modes ϕ_I and the gaugino zero modes λ^a:

$$-\int d^4x \left(\frac{1}{2}m^2_{IJ}\phi^*_I\phi_J + \frac{1}{6}A_{IJK}\phi_I\phi_J\phi_K + \frac{1}{2}M_a\lambda^a\lambda^a + \text{h.c.} \right). \tag{22}$$

Under the SS twist (9), the gaugino zero mode receives a soft mass

$$M_a = -\frac{\omega}{R}.$$

From (16) and (18), one easily finds the KK wavefunctions of the hyperscalar zero modes:

$$\begin{aligned} \tilde{\phi}^1_I &= \left(\frac{M_I}{1-e^{-2M_I\pi R}}\right)^{1/2} e^{-M_I R y} + \mathcal{O}(\omega^2), \\ \tilde{\phi}^2_I &= \pi\omega\left(\frac{M_I}{e^{2M_I\pi R}-1}\right)^{1/2} \frac{e^{M_I R y}-e^{-M_I R y}}{e^{M_I\pi R}-e^{-M_I\pi R}} + \mathcal{O}(\omega^3), \end{aligned} \tag{23}$$

and their soft masses:

$$m^2_{IJ} = m^2_I\delta_{IJ} = \left(\frac{\omega}{R}\right)^2 \left(\frac{M_I\pi R}{\sinh(M_I\pi R)}\right)^2 + \mathcal{O}(\omega^4). \tag{24}$$

The Yukawa couplings and A-parameters of these zero modes can be obtained from (20) and (21), yielding

$$\begin{aligned} y_{IJK} &= \frac{1}{\sqrt{Y_IY_JY_K}}\Big(\lambda_{IJK} + \lambda'_{IJK}e^{-(M_I+M_J+M_K)\pi R} \Big) + \mathcal{O}(\omega^2), \\ A_{IJK} &= \frac{\omega}{R}\frac{1}{\sqrt{Y_IY_JY_K}}\left[\left(\frac{2M_I\pi R}{e^{2M_I\pi R}-1} + \frac{2M_J\pi R}{e^{2M_J\pi R}-1} + \frac{2M_K\pi R}{e^{2M_K\pi R}-1} \right)\lambda_{IJK} \right. \\ &+ \left. \left(\frac{2M_I\pi R}{1-e^{-2M_I\pi R}} + \frac{2M_J\pi R}{1-e^{-2M_J\pi R}} + \frac{2M_K\pi R}{1-e^{-2M_K\pi R}} \right)\lambda'_{IJK}e^{-(M_I+M_J+M_K)\pi R} \right] + \mathcal{O}(\omega^3)\,, \end{aligned} \tag{25}$$

where

$$Y_I = \frac{1}{M_I}\left(1 - e^{-2M_I\pi R}\right).$$

Note that the above Yukawa couplings and soft parameters of zero modes should be interpreted as the parameters renormalized at the compactification scale $M_c \sim 1/R$.

It has been pointed out that the SS SUSY breaking in 5D orbifold SUGRA has an interesting correspondence with the radion-mediated SUSY breaking. Here we explicitly show that the zero mode soft parameters induced by the SS boundary condition are precisely same as the radion-mediated soft parameters in 4D effective theory. To see this, let us construct the 4D effective action of the gauge and matter zero modes without any SS twist, i.e. $\omega = 0$, while keeping the radion superfield T to take a generic value. The resulting 4D effective action can be written on $N = 1$ superspace, and can be obtained easily by making the radion-dependent superfield redefinition:

$$H_I \to e^{M_I T|y|} H_I, \quad H_I^c \to e^{-M_I T|y|} H_I^c.$$

After this field redefinition, the bulk and boundary actions of (10) and (13) become

$$\begin{aligned} S_{\rm bulk} &= \int d^5x \left[\int d^4\theta \, \frac{T+T^*}{2} \left(e^{-M_I(T+T^*)|y|} H_I H_I^* + e^{M_I(T+T^*)|y|} H_I^c H_I^{c*} \right) \right. \\ &\qquad \left. + \int d^2\theta \frac{1}{4g_{5a}^2} T W^{a\alpha} W^a_\alpha + \text{h.c.} \right], \\ S_{\rm brane} &= \int d^5x \int d^2\theta \left(\delta(y) \frac{1}{6} \lambda_{IJK} H_I H_J H_K + \delta(y-\pi) \frac{1}{6} \lambda'_{IJK} e^{-(M_I+M_J+M_K)T|y|} H_I H_J H_K \right) + \text{h.c.} \end{aligned} \tag{26}$$

In the new 5D superfield basis, all zero modes have *constant* wavefunctions, thus their 4D effective action can be obtained by simply integrating the 5D action over the 5-th dimension. Let Φ_I denote the constant zero modes of H_I, and W^a_α denote the field strength superfields for the constant zero modes of V^a. We then find

$$S_{4D} = \int d^4x \left[\int d^4\theta \, Y_{I\bar{J}} \Phi_I \Phi_J^* + \int d^2\theta \left(\frac{1}{4} f_a W^{a\alpha} W^a_\alpha + \tilde{y}_{IJK} \Phi_I \Phi_J \Phi_K \right) + \text{h.c.} \right], \tag{27}$$

where the hermitian wavefunction coefficients $Y_{I\bar{J}}$, the holomorphic Yukawa couplings $\tilde{y}_{IJK}$, and the holomorphic gauge kinetic functions f_a are given by

$$\begin{aligned} Y_{I\bar{J}} &= Y_I \delta_{IJ} = \frac{1}{M_I}\left(1 - e^{-\pi M_I(T+T^*)}\right) \delta_{IJ}, \\ \tilde{y}_{IJK} &= \lambda_{IJK} + \lambda'_{IJK} e^{-\pi(M_I+M_J+M_K)T}, \\ f_a &= \frac{2\pi}{g_{5a}^2} T. \end{aligned} \tag{28}$$

If the radion F-component, F^T, is the major source of SUSY breaking in the above 4D effective action, the soft parameters of the canonically normalized 4D fields are given by

$$\begin{aligned} M_a &= -\frac{1}{2\,\text{Re}(f_a)} F^T \partial_T f_a = -\frac{F^T}{2R}, \\ m^2_{I\bar{J}} &= -\frac{1}{\sqrt{Y_I Y_J}} \left|F^T\right|^2 \left(\partial_T \partial_{T^*} Y_{I\bar{J}} - \Gamma_{T\,I}^{\,K} \partial_{T^*} Y_{K\bar{J}} \right) \\ &= \delta_{IJ} \left| \frac{F^T}{2R} \right|^2 \left(\frac{M_I \pi R}{\sinh(M_I \pi R)} \right)^2, \\ A_{IJK} &= -\frac{1}{\sqrt{Y_I Y_J Y_K}} F^T \left(\partial_T \tilde{y}_{IJK} - \Gamma_{T\,I}^{\,L} \tilde{y}_{LJK} - \Gamma_{T\,J}^{\,L} \tilde{y}_{ILK} - \Gamma_{T\,K}^{\,L} \tilde{y}_{IJL} \right) \\ &= \frac{F^T}{2R} \frac{1}{\sqrt{Y_I Y_J Y_K}} \left[\left(\frac{2M_I\pi R}{e^{2M_I\pi R}-1} + \frac{2M_J\pi R}{e^{2M_J\pi R}-1} + \frac{2M_K\pi R}{e^{2M_K\pi R}-1} \right) \lambda_{IJK} \right. \\ &\quad \left. + \left(\frac{2M_I\pi R}{1-e^{-2M_I\pi R}} + \frac{2M_J\pi R}{1-e^{-2M_J\pi R}} + \frac{2M_K\pi R}{1-e^{-2M_K\pi R}} \right) \lambda'_{IJK} e^{-(M_I+M_J+M_K)\pi R} \right], \end{aligned} \tag{29}$$

where the Kähler connection $\Gamma_T{}^I_J = Y^{I\bar{K}}\partial_T Y_{J\bar{K}}$. Obviously the above radion-mediated soft parameters are precisely same as the SS-induced soft parameters (up to small corrections higher order in ω) with the matching condition

$$F^T = 2\omega\,.$$

This means that our phenomenological analysis of SS-induced soft parameters in the next section can be applied to any SUSY breaking mechanism giving a sizable F^T, for instance the hidden gaugino condensation model.

So far, we have considered the most general scenario that the Yukawa couplings originate from both fixed points, $y = 0$ and π. In fact, to generate hierarchical Yukawa couplings through quasi-localization in a natural manner, one needs to assume that Yukawa couplings originate *only* from one fixed point, e.g. from $y = 0$. In this case, the Yukawa couplings and soft parameters at the compactification scale are given by

$$\begin{aligned}
y_{IJK} &= \left(\frac{M_I M_J M_K}{(1-e^{-2M_I\pi R})(1-e^{-2M_J\pi R})(1-e^{-2M_K\pi R})}\right)^{\frac{1}{2}}\lambda_{IJK} \\
A_{IJK} &= \frac{\omega}{R}\left(\frac{2M_I\pi R}{e^{2M_I\pi R}-1}+\frac{2M_J\pi R}{e^{2M_J\pi R}-1}+\frac{2M_K\pi R}{e^{2M_K\pi R}-1}\right)y_{IJK} \\
M_a &= -\frac{\omega}{R}\,, \\
m_I^2 &= \left(\frac{\omega}{R}\right)^2\left(\frac{M_I\pi R}{\sinh(M_I\pi R)}\right)^2.
\end{aligned} \tag{30}$$

As we have noticed, the matter zero mode Φ_I from a 5D hypermultiplet with kink mass M_I has an wavefunction of the form $e^{-M_I R|y|}$, thus is quasi-localized at $y = 0$ ($y = \pi$) if $M_I > 0$ ($M_I < 0$). As a result, in case that Yukawa couplings originate from $y = 0$, the quark/lepton superfields from hypermultiplets with $M_I < 0$ would have (exponentially) small Yukawa couplings, while the quark/lepton superfields with $M_I > 0$ can have Yukawa couplings of order unity. Obviously, the above form of Yukawa couplings shows this feature, achieving the Yukawa hierarchy from quasi-localization. The above results show also that the soft scalar masses m_I^2 and the A to Yukawa ratios A_{IJK}/y_{IJK} are highly flavor-dependent at the compactification scale. Although the flavor-violating pieces are suppressed with an appropriate correlation with Yukawa couplings, still they can lead to dangerous flavor-violations at low energy scales as will be discussed in the next section.

Assuming that the Higgs superfields are boundary superfields confined at $y = 0$ simplifies the form of Yukawa and A-parameters, however their flavor structures are essentially the same. In our framework, any boundary superfield can be interpreted as the zero mode of bulk hypermultiplet having an infinite kink mass, more precisely a kink mass comparable to the 5D cutoff scale Λ_5. Note that in this case all other KK modes have the masses of $\mathcal{O}(\Lambda_5)$, so are decoupled. Then the Yukawa and A-parameters of boundary Higgs superfields Φ_K can be obtained from (30) by taking the limit $M_K \to \Lambda_5$ together with an appropriate redefinition of boundary Yukawa couplings, yielding

$$\begin{aligned}
y_{IJ} &= \left[\frac{M_I M_J}{(1-e^{-2M_I\pi R})(1-e^{-2M_J\pi R})}\right]^{\frac{1}{2}}\tilde{\lambda}_{IJ} \\
A_{IJ} &= \frac{\omega}{R}\left(\frac{2M_I\pi R}{e^{2M_I\pi R}-1}+\frac{2M_J\pi R}{e^{2M_I\pi R}-1}\right)y_{IJ}
\end{aligned} \tag{31}$$

where

$$\tilde{\lambda}_{IJ} \equiv \left(\frac{M_K}{1-e^{-2M_K\pi R}}\right)^{1/2}\lambda_{IJK} \approx \Lambda_5^{1/2}\lambda_{IJK} \tag{32}$$

for the quark/lepton flavor indices I, J.

In the next section, we analyze in detail the resulting low energy flavor violations under the assumption that the kink masses M_I are appropriately quantized. Note that in 5D orbifold SUGRA the assumption of quantized kink masses corresponds to the assumption of quantized $U(1)$ gauge charges. It is then convenient to write the above Yukawa couplings and soft parameters in the following way:

$$\begin{aligned}
y_{IJ} &= \tilde{\lambda}_{IJ}\frac{\ln(1/\epsilon)}{\pi R}\sqrt{\frac{N_I N_J}{(\epsilon^{-2N_I}-1)(\epsilon^{-2N_J}-1)}}\,, \\
M_a &= -\frac{\omega}{R}\,,
\end{aligned}$$

$$
\begin{aligned}
A_{IJ} &= 2y_{IJ}\ln(1/\epsilon)\frac{\omega}{R}\left(\frac{N_I}{1-\epsilon^{2N_I}}+\frac{N_J}{1-\epsilon^{2N_J}}\right), \\
m^2_{I\bar{J}} &= \delta_{IJ}\left(2\ln(1/\epsilon)\frac{N_I}{\epsilon^{N_I}-\epsilon^{-N_I}}\frac{\omega}{R}\right)^2,
\end{aligned}
\tag{33}
$$

where

$$
N_I = -\frac{\pi R}{\ln(1/\epsilon)}M_I \quad \text{for} \quad \epsilon \equiv \text{Cabibbo angle} \approx 0.2.
$$

The above Yukawa couplings are quite similar to the Yukawa couplings in Frogatt-Nielsen models with N_I being identified as the $U(1)_F$ charges. More explicitly,

$$
y_{IJ} \simeq \frac{\tilde{\lambda}_{IJ}\ln(1/\epsilon)}{\pi R}\sqrt{Z_I Z_J}\,\epsilon^{X_I+X_J} \equiv \lambda_{IJ}\epsilon^{X_I+X_J}\,, \tag{34}
$$

where

$$
Z_I = \begin{cases} |N_I| & (\,N_I \neq 0\,) \\ 1/[2\ln(1/\epsilon)] & (\,N_I = 0\,)\,, \end{cases} \tag{35}
$$

and the effective flavor charge X_I is given by

$$
X_I = \begin{cases} N_I & (\,N_I \geq 0\,) \\ 0 & (\,N_I < 0\,)\,. \end{cases} \tag{36}
$$

When the theory is strongly coupled at Λ_5, a naive dimensional analysis [15] suggests that

$$
\Lambda_5 \pi R = \mathcal{O}(6\pi^3), \quad \tilde{\lambda}_{IJ} = \mathcal{O}(\sqrt{6\pi^3}/\Lambda_5).
$$

Then the redefined boundary Yukawa couplings $\lambda_{IJ} \simeq \tilde{\lambda}_{IJ}\sqrt{M_I M_J}$ would be of order unity if the corresponding kink masses $|M_I| = |N_I|\ln(1/\epsilon)/\pi R = \mathcal{O}(\sqrt{6\pi^3}/\pi R)$. In the following, we will ignore the factor $2 \sim 3$ differences of λ_{IJ} arising from their M_I-dependence, and simply assume that the redefined boundary Yukawa couplngs λ_{IJ} are all of order unity. Then the observed quark/lepton masses and CKM mixing angles can be explained by the 5D kink masses M_I which are quantized in a manner to give integer-valued N_I.

As for the soft scalar masses and trilinear couplings, the above results can be approximated as

$$
\begin{aligned}
A_{IJ} &\simeq M_0\, y_{IJ}\,(a_I + a_J)\,, \\
m^2_{I\bar{J}} &\simeq \delta_{I\bar{J}} M_0^2 \begin{cases} N_I^2 \epsilon^{2|N_I|} & (\,N_I \neq 0\,) \\ 1/[2\ln(1/\epsilon)]^2 & (\,N_I = 0\,) \end{cases}
\end{aligned}
\tag{37}
$$

where

$$
M_0 = 2\ln(1/\epsilon)\frac{\omega}{R}, \tag{38}
$$

and

$$
a_I = \begin{cases} N_I & (\,N_I > 0\,) \\ 1/[2\ln(1/\epsilon)] & (\,N_I = 0\,) \\ |N_I|\epsilon^{2|N_I|} & (\,N_I < 0\,)\,. \end{cases} \tag{39}
$$

An important feature of the SS SUSY breaking is that A_{IJ}/y_{IJ} and $m^2_{I\bar{J}}$ are *quantized* for $N_I \neq 0$ in the leading approximation. This feature of the SS SUSY breaking, more generally of the radion-mediated SUSY breaking, is quite useful for suppressing dangerous flavor violations. With this feature, flavor violating amplitudes appear in a form $f(N_I) - f(N_J)$, thus are canceled if some of N_I are degenerate.

The suppression of $m^2_{I\bar{J}}/M_0^2$ and A_{IJ}/M_0 by some powers of ϵ is essentially due to the quasi-localization of matter zero modes. The SUSY breaking by boundary condition is a non-local SUSY breaking, so the resulting soft parameters are more suppressed for more localized matter fields. Note that the suppressions of y_{IJ} and A_{IJ} are asymmetric under $N_I \to -N_I$. This is simply because the Yukawa couplings originate from the boundary at $y = 0$. On the other hand, $m^2_{I\bar{J}}$ are independent of the origin of the Yukawa couplings, so are symmetric under $N_I \to -N_I$.

III. LOW ENERGY FLAVOR VIOLATIONS

In this section, we analyze the low energy flavor violations resulting from the SS SUSY breaking for quasi-localized matter fields. The renormalization scale for the SS-induced soft parameters of (33) can be identified as the compactification scale M_c. To be specific, here we assume that M_c is close to the unification scale $M_{GUT} \sim 2 \times 10^{16}$ GeV and the 4D effective theory below M_c is given by the MSSM. For simplicity, we further assume that the two MSSM Higgs doublets H_1 and H_2 are boundary superfields confined at $y = 0$.

Let $\psi_I = \{q_i, u_i, d_i, \ell_i, e_i\}$ $(i = 1, 2, 3)$ denote the known three generations of the left-handed quark-doublets (q_i), up-type antiquark-singlets (u_i), down-type antiquark singlets (d_i), lepton-doublets (ℓ_i), and anti-lepton singlets (e_i). The Yukawa couplings can be written as

$$\mathcal{L}_{\rm Yukawa} = y^u_{ij} H_2 u_i q_j + y^d_{ij} H_1 d_i q_j + y^\ell_{ij} H_1 e_i \ell_j \tag{40}$$

and the squark/sleptons $\phi_I = \{\tilde{q}_i, \tilde{u}_i, \tilde{d}_i, \tilde{\ell}_i, \tilde{e}_i\}$ have the soft SUSY breaking couplings:

$$\begin{aligned}\mathcal{L}_{\rm soft} \;=\; -\Big(& A^u_{ij} H_2 \tilde{u}_i \tilde{q}_j + A^d_{ij} H_1 \tilde{d}_i \tilde{q}_j + A^\ell_{ij} H_1 \tilde{e}_i \tilde{\ell}_j \\ & + m^{2(\tilde{q})}_{i\bar{j}} \tilde{q}_i \tilde{q}^*_j + m^{2(\tilde{u})}_{i\bar{j}} \tilde{u}_i \tilde{u}^*_j + m^{2(\tilde{d})}_{i\bar{j}} \tilde{d}_i \tilde{d}^*_j \\ & + m^{2(\tilde{\ell})}_{i\bar{j}} \tilde{\ell}_i \tilde{\ell}^*_j + m^{2(\tilde{e})}_{i\bar{j}} \tilde{e}_i \tilde{e}^*_j \Big) .\end{aligned} \tag{41}$$

As shown in (34), the canonical 4D Yukawa couplings in (40) are given by

$$y^u_{ij} \simeq \lambda^u_{ij} \epsilon^{X^u_i + X^q_j}, \quad y^d_{ij} \simeq \lambda^d_{ij} \epsilon^{X^d_i + X^q_j}, \quad y^\ell_{ij} \simeq \lambda^\ell_{ij} \epsilon^{X^e_i + X^\ell_j}, \tag{42}$$

where the redefined boundary couplings λ^ψ_{ij} $(\psi = u, d, \ell)$ are given by

$$\lambda^\psi_{ij} = \frac{\tilde{\lambda}^\psi_{ij} \ln(1/\epsilon)}{\pi R} \sqrt{Z^\psi_i Z^\psi_j}. \tag{43}$$

Here $\tilde{\lambda}^\psi_{ij}$ are the boundary Yukawa couplings which are generically of order $\sqrt{6\pi^3}/\Lambda_5$ [15],

$$Z^\psi_i = \begin{cases} |N^\psi_i| & (\, N^\psi_i \neq 0 \,) \\ 1/[2\ln(1/\epsilon)] & (\, N^\psi_i = 0 \,), \end{cases} \tag{44}$$

and X^ψ_i $(\psi = q, u, d, \ell)$ are the effective flavor charges defined as

$$X^\psi_i = \begin{cases} N^\psi_i & (\, N^\psi_i \geq 0 \,) \\ 0 & (\, N^\psi_i < 0 \,). \end{cases} \tag{45}$$

In the following, we assume that N^ψ_i are all integers and X^ψ_i take the normal hierarchy as

$$X^\psi_1 \geq X^\psi_2 \geq X^\psi_3 .$$

Neglecting the part of A_{IJ}/y_{IJ} suppressed by $\epsilon^{2|N_I|}$ or $\epsilon^{2|N_J|}$ in (37), the soft parameters of (41) can be approximated as

$$\begin{aligned} A^u_{ij} &\simeq M_0 (X^u_i + X^q_j)\, y^u_{ij}, \\ A^d_{ij} &\simeq M_0 (X^d_i + X^q_j)\, y^d_{ij}, \\ A^\ell_{ij} &\simeq M_0 (X^e_i + X^\ell_j)\, y^\ell_{ij}, \\ m^{2(\tilde{\psi})}_{i\bar{j}} &\simeq \delta_{i\bar{j}} M_0^2 \begin{cases} |N^\psi_i|^2 \epsilon^{2|N^\psi_i|} & (\, N^\psi_i \neq 0 \,) \\ 1/[2\ln(1/\epsilon)]^2 & (\, N^\psi_i = 0 \,), \end{cases} \end{aligned} \tag{46}$$

where

$$M_0 = 2\ln(1/\epsilon) \left(\frac{\omega}{R} \right) = -2 M_{1/2} \ln(1/\epsilon)$$

for the universal gaugino mass $M_{1/2} = -\omega/R$. The above flavor structure of trilinear A-couplings is shared by various SUSY–breaking scenarios [16], while that of soft scalar masses is rather specific to the SS SUSY breaking.

To perform the analysis of low energy observables, let us first introduce a parameterization of λ^ψ_{ij} ($\psi = u, d, \ell$) which are assumed to be of order unity. As an example, the down-type Yukawa matrix at the unification scale can be decomposed as

$$y^d_{ij} \;=\; V_d^\dagger \begin{pmatrix} \hat{y}^d_1 & & \\ & \hat{y}^d_2 & \\ & & \hat{y}^d_3 \end{pmatrix} V_q\,, \tag{47}$$

where V_d, V_q are 3×3 unitary matrices and $\hat{y}^d_i$ can be determined from the observed quark masses. With this expression, λ^d_{ij} can be written as

$$\begin{aligned} \lambda^d_{ij} \;=\;& (\hat{y}^d_1 \epsilon^{-X^d_1 - X^q_1})\,(V^*_d)_{1i}\frac{\epsilon^{X^d_1}}{\epsilon^{X^d_i}}\,(V_q)_{1j}\frac{\epsilon^{X^q_1}}{\epsilon^{X^q_j}} + (\hat{y}^d_2 \epsilon^{-X^d_2 - X^q_2})\,(V^*_d)_{2i}\frac{\epsilon^{X^d_2}}{\epsilon^{X^d_i}}\,(V_q)_{2j}\frac{\epsilon^{X^q_2}}{\epsilon^{X^q_j}} \\ &+(\hat{y}^d_3 \epsilon^{-X^d_3 - X^q_3})\,(V^*_d)_{3i}\frac{\epsilon^{X^d_3}}{\epsilon^{X^d_i}}\,(V_q)_{3j}\frac{\epsilon^{X^q_3}}{\epsilon^{X^q_j}}. \end{aligned} \tag{48}$$

Barring an accidental cancellation between different terms, each term in (48) should not exceed $\mathcal{O}(1)$. Noting that $\hat{y}^d_i = \mathcal{O}(\epsilon^{X^d_i + X^q_i})$ and also using the unitarity of the mixing matrices, we find the following order of magnitude constraints from the second and third terms of (48):

$$\begin{aligned} |(V_{d,q})_{21}|\sqrt{1-|(V_{q,d})_{23}|^2} &\lesssim \frac{\epsilon^{X^{d,q}_1}}{\epsilon^{X^{d,q}_2}}, \\ |(V_{d,q})_{31}|\sqrt{1-|(V_{q,d})_{23}|^2} &\lesssim \frac{\epsilon^{X^{d,q}_1}}{\epsilon^{X^{d,q}_3}}, \\ |(V_{d,q})_{32}|\sqrt{1-|(V_{q,d})_{13}|^2} &\lesssim \frac{\epsilon^{X^{d,q}_2}}{\epsilon^{X^{d,q}_3}}. \end{aligned} \tag{49}$$

Except for the special crossing points with $|(V_{q,d})_{13}| \simeq 1$ or $|(V_{q,d})_{23}| \simeq 1$, these constraints are reduced to

$$|(V_{d,q})_{21}| \lesssim \frac{\epsilon^{X^{d,q}_1}}{\epsilon^{X^{d,q}_2}}, \quad |(V_{d,q})_{31}| \lesssim \frac{\epsilon^{X^{d,q}_1}}{\epsilon^{X^{d,q}_3}}, \quad |(V_{d,q})_{32}| \lesssim \frac{\epsilon^{X^{d,q}_2}}{\epsilon^{X^{d,q}_3}}. \tag{50}$$

In general, the unitary matrix V_d can be decomposed as

$$V_d \;=\; e^{i\varphi_d} e^{i\vec{\phi}^T_d}\,\overline{V}_d\, e^{i\vec{\psi}_d}, \quad (\sum_i \phi^i_d = \sum_i \psi^i_d = 0)\,, \tag{51}$$

where

$$\begin{aligned} (\overline{V}_d)_{ij} \;&=\; \begin{pmatrix} c_{12} & -s_{12} & 0 \\ s_{12} & c_{12} & 0 \\ 0 & 0 & 1 \end{pmatrix} \begin{pmatrix} c_{13} & 0 & -s_{13}e^{-i\delta^D_{13}} \\ 0 & 1 & 0 \\ s_{13}e^{i\delta^D_{13}} & 0 & c_{13} \end{pmatrix} \begin{pmatrix} 1 & 0 & 0 \\ 0 & c_{23} & -s_{23} \\ 0 & s_{23} & c_{23} \end{pmatrix} \\ &=\; \begin{pmatrix} c_{12}c_{13} & -s_{12}c_{23} - c_{12}s_{23}s_{13}e^{-i\delta^D_{13}} & s_{12}s_{23} - c_{12}c_{23}s_{13}e^{-i\delta^D_{13}} \\ s_{12}c_{13} & c_{12}c_{23} - s_{12}s_{23}s_{13}e^{-i\delta^D_{13}} & -c_{12}s_{23} - s_{12}c_{23}s_{13}e^{-i\delta^D_{13}} \\ s_{13}e^{i\delta^D_{13}} & s_{23}c_{13} & c_{23}c_{13} \end{pmatrix}, \end{aligned} \tag{52}$$

for $s_{ij} \equiv \sin\theta^d_{ij}$ and $c_{ij} \equiv \cos\theta^d_{ij}$. With this parameterization, the order of magnitude constraints (50) are translated into

$$\theta^d_{12} \lesssim \frac{\epsilon^{X^d_1}}{\epsilon^{X^d_2}}, \quad \theta^d_{13} \lesssim \frac{\epsilon^{X^d_1}}{\epsilon^{X^d_3}}, \quad \theta^d_{23} \lesssim \frac{\epsilon^{X^d_2}}{\epsilon^{X^d_3}}. \tag{53}$$

Together with the expression (51) and also a similar expression of V_q, the above constraints imply

$$|(V_{d,q})_{12}| \lesssim \frac{\epsilon^{X^{d,q}_1}}{\epsilon^{X^{d,q}_2}}, \quad |(V_{d,q})_{13}| \lesssim \frac{\epsilon^{X^{d,q}_1}}{\epsilon^{X^{d,q}_3}}, \quad |(V_{d,q})_{23}| \lesssim \frac{\epsilon^{X^{d,q}_2}}{\epsilon^{X^{d,q}_3}}. \tag{54}$$

For later discussion, it is convenient to introduce the following $\mathcal{O}(1)$ parameters:

$$(\kappa_{d,q})_{ij} = \frac{(\overline{V}_{d,q})_{ij}}{\epsilon^{|X_i^{d,q} - X_j^{d,q}|}}. \tag{55}$$

Then using the freedom of rephasing the fields, we finally arrive at the following parameterizations and also the constraints, reflecting well the underlying structure (42):

$$\begin{aligned}
y^u_{ij} &= \overline{V}^\dagger_u \begin{pmatrix} \hat{y}^u_1 e^{i\phi^u_1} & & \\ & \hat{y}^u_2 e^{i\phi^u_2} & \\ & & \hat{y}^u_3 e^{i\phi^u_3} \end{pmatrix} V_{CKM} \overline{V}_q, \\
y^d_{ij} &= \overline{V}^\dagger_d \begin{pmatrix} \hat{y}^d_1 e^{i\phi^d_1} & & \\ & \hat{y}^d_2 e^{i\phi^d_2} & \\ & & \hat{y}^d_3 e^{i\phi^d_3} \end{pmatrix} \overline{V}_q, \\
y^\ell_{ij} &= \overline{V}^\dagger_e \begin{pmatrix} \hat{y}^\ell_1 e^{i\phi^\ell_1} & & \\ & \hat{y}^\ell_2 e^{i\phi^\ell_2} & \\ & & \hat{y}^\ell_3 e^{i\phi^\ell_3} \end{pmatrix} \overline{V}_\ell,
\end{aligned} \tag{56}$$

$$\begin{aligned}
&\theta^\psi_{12} \lesssim \frac{\epsilon^{X^\psi_1}}{\epsilon^{X^\psi_2}}, \quad \theta^\psi_{13} \lesssim \frac{\epsilon^{X^\psi_1}}{\epsilon^{X^\psi_3}}, \quad \theta^\psi_{23} \lesssim \frac{\epsilon^{X^\psi_2}}{\epsilon^{X^\psi_3}}, \\
&(\kappa_\psi)_{ij} \equiv \frac{(\overline{V}_\psi)_{ij}}{\epsilon^{|X^\psi_i - X^\psi_j|}} \simeq \mathcal{O}(1),
\end{aligned} \tag{57}$$

where $\phi^\psi_1 + \phi^\psi_2 + \phi^\psi_3 = 0$ for $\psi = q, u, d, \ell, e$.

Assuming that the CKM matrix does not involve any fine-tuned cancellation among mixing angles, we obtain

$$\theta^q_{12} \simeq \epsilon, \quad \theta^q_{23} \simeq \epsilon^2, \quad \theta^q_{13} \simeq \epsilon^3. \tag{58}$$

Also the observed fermion mass spectrum indicates

$$\begin{aligned}
&\frac{\hat{y}^u_1}{\hat{y}^u_3} \simeq 7.2 \times 10^{-6} \simeq \epsilon^{7-8}, \quad \frac{\hat{y}^u_2}{\hat{y}^u_3} \simeq 2.4 \times 10^{-3} \simeq \epsilon^{3-4}, \\
&\frac{\hat{y}^d_1}{\hat{y}^d_3} \simeq 1.0 \times 10^{-3} \simeq \epsilon^{4-5}, \quad \frac{\hat{y}^d_2}{\hat{y}^d_3} \simeq 2.0 \times 10^{-2} \simeq \epsilon^{2-3}, \\
&\frac{\hat{y}^\ell_1}{\hat{y}^\ell_3} \simeq 2.8 \times 10^{-4} \simeq \epsilon^{5-6}, \quad \frac{\hat{y}^\ell_2}{\hat{y}^\ell_3} \simeq 5.9 \times 10^{-2} \simeq \epsilon^{1-2},
\end{aligned} \tag{59}$$

where we assume that $\tan\beta = \langle H^0_2 \rangle / \langle H^0_1 \rangle$ is not so large. In fact, as for the flavor conserving part, the SS SUSY breaking with quasi-localized matter fields is somewhat similar to the gaugino-mediation [17] or the no-scale model since the soft parameters of matter fields at the compactification scale are suppressed by the quasi-localization factor. As a result, the model gives either a stau LSP or a negative stau mass-square when $\tan\beta$ is large [18]. Together with this consideration, X^ψ_i favored by the observed fermion masses and CKM mixing angles are given by [19]

$$\begin{aligned}
&\vec{X}^q = (3, 2, 0), \quad \vec{X}^u = (4 \text{ or } 5,\ 1 \text{ or } 2,\ 0), \\
&\vec{X}^d = (1 + x \text{ or } 2 + x,\ x \text{ or } 1 + x,\ x), \\
&\vec{X}^\ell + \vec{X}^e = (6 + x \text{ or } 5 + x,\ 2 + x \text{ or } 1 + x,\ x),
\end{aligned} \tag{60}$$

where $x = 1, 2$ corresponds to medium and small $\tan\beta$, respectively.

We are now ready to discuss the phenomenology of flavor mixings resulting from the SS-mediated soft parameters (46). To this end, it is convenient to rotate the fields to the super-CKM basis in which the Yukawa couplings have diagonal form up to V_{CKM}. In the SCKM basis, the Higgs-slepton trilinear couplings are given by

$$(V_e A^\ell V^\dagger_\ell)_{ij} \simeq M_0 \left[e^{i(\phi^\ell_i - \phi^\ell_j)} (\overline{V}_e)_{ik} X^e_k (\overline{V}^\dagger_e)_{kj}\, \hat{y}^\ell_j + \hat{y}^\ell_i\, (\overline{V}_\ell)_{ik} X^\ell_k (\overline{V}^\dagger_\ell)_{kj} \right]. \tag{61}$$

Using the unitarity of $V_{e,\ell}$, one easily finds

$$
\begin{aligned}
(\overline{V}_e)_{ik} X^e_k (\overline{V}^\dagger_e)_{kj} &= (\overline{V}_e)_{i1}(\overline{V}^*_e)_{j1}(X^e_1 - X^e_2) - (\overline{V}_e)_{i3}(\overline{V}^*_e)_{j3}(X^e_2 - X^e_3) + \delta_{ij} X^e_2 \\
&= -(\overline{V}_e)_{i2}(\overline{V}^*_e)_{j2}(X^e_1 - X^e_2) - (\overline{V}_e)_{i3}(\overline{V}^*_e)_{j3}(X^e_1 - X^e_3) + \delta_{ij} X^e_1
\end{aligned}
\tag{62}
$$

and also a similar relation for X^ℓ and $\overline{V}_\ell$. This shows that flavor violation is more suppressed in case that the lepton superfields in different generation have a common effective flavor charge.

To parameterize the flavor mixing, let us introduce the $\tilde{\delta}$ parameters *at the compactification scale*. The $\tilde{\delta}$ parameters for RL/LR mixing are defined as

$$
(\tilde{\delta}^\ell_{RL})_{ij} = (\tilde{\delta}^\ell_{LR})^*_{ji} \equiv \frac{(V_e A^\ell V^\dagger_\ell)_{ij}}{|M_{1/2}|^2} v_d,
\tag{63}
$$

where $v_d \simeq 174\text{GeV}\cos\beta = \langle H^0_1 \rangle$. Using the unitarity relations, we then find

$$
\begin{aligned}
(\tilde{\delta}^\ell_{RL})_{12} &\simeq -6.44\times 10^{-4}\left(\frac{500\text{GeV}}{|M_{1/2}|}\right)\left(\frac{m_\mu(M_{GUT})}{100\text{MeV}}\right) \\
&\quad \times \left[e^{i(\phi^\ell_1-\phi^\ell_2)} \epsilon^{\Delta X^e_{12}} \left\{ -\Delta X^e_{12} + \epsilon^{2\Delta X^e_{23}}\ \Delta X^e_{23}\ \frac{(\kappa_e)_{13}(\kappa^*_e)_{23}}{(\kappa_e)_{11}(\kappa^*_e)_{21}} \right\} (\kappa_e)_{11}(\kappa^*_e)_{21} \right. \\
&\quad \left. + \left(\frac{m_e}{m_\mu}\right) \epsilon^{\Delta X^\ell_{12}} \left\{ -\Delta X^\ell_{12} + \epsilon^{2\Delta X^\ell_{23}}\ \Delta X^\ell_{23}\ \frac{(\kappa_\ell)_{13}(\kappa^*_\ell)_{23}}{(\kappa_\ell)_{11}(\kappa^*_\ell)_{21}} \right\} (\kappa_\ell)_{11}(\kappa^*_\ell)_{21} \right], \\
(\tilde{\delta}^\ell_{RL})_{13} &\simeq -6.44\times 10^{-3}\left(\frac{500\text{GeV}}{|M_{1/2}|}\right)\left(\frac{m_\tau(M_{GUT})}{1.00\text{GeV}}\right) \\
&\quad \times \left[e^{i(\phi^\ell_1-\phi^\ell_3)} \epsilon^{\Delta X^e_{13}} \left\{ \Delta X^e_{13} + \frac{(\kappa_e)_{12}(\kappa^*_e)_{32}}{(\kappa_e)_{13}(\kappa^*_e)_{33}} \Delta X^e_{12} \right\} (\kappa_e)_{13}(\kappa^*_e)_{33} \right. \\
&\quad \left. + \left(\frac{m_e}{m_\tau}\right) \epsilon^{\Delta X^\ell_{13}} \left\{ \Delta X^\ell_{13} + \frac{(\kappa_\ell)_{12}(\kappa^*_\ell)_{32}}{(\kappa_\ell)_{13}(\kappa^*_\ell)_{33}} \Delta X^\ell_{12} \right\} (\kappa_\ell)_{13}(\kappa^*_\ell)_{33} \right], \\
(\tilde{\delta}^\ell_{RL})_{23} &\simeq -6.44\times 10^{-3}\left(\frac{500\text{GeV}}{|M_{1/2}|}\right)\left(\frac{m_\tau(M_{GUT})}{1.00\text{GeV}}\right) \\
&\quad \times \left[e^{i(\phi^\ell_2-\phi^\ell_3)} \epsilon^{\Delta X^e_{23}} \left\{ \Delta X^e_{23} - \epsilon^{2\Delta X^e_{12}} \frac{(\kappa_e)_{21}(\kappa^*_e)_{31}}{(\kappa_e)_{23}(\kappa^*_e)_{33}} \Delta X^e_{12} \right\} (\kappa_e)_{23}(\kappa^*_e)_{33} \right. \\
&\quad \left. + \left(\frac{m_\mu}{m_\tau}\right) \epsilon^{\Delta X^\ell_{23}} \left\{ \Delta X^\ell_{23} - \epsilon^{2\Delta X^\ell_{12}} \frac{(\kappa_\ell)_{21}(\kappa^*_\ell)_{31}}{(\kappa_\ell)_{23}(\kappa^*_\ell)_{33}} \Delta X^\ell_{12} \right\} (\kappa_\ell)_{23}(\kappa^*_\ell)_{33} \right],
\end{aligned}
\tag{64}
$$

where $(\kappa_\psi)_{ij} = \mathcal{O}(1)$ as defined in (57) and $\Delta X^\psi_{ij} \equiv X^\psi_i - X^\psi_j$.

Note that if $X^e_1 = X^e_2$ or $X^e_2 = X^e_3$, $(\tilde{\delta}^\ell_{RL})_{12}$ or $(\tilde{\delta}^\ell_{RL})_{23}$ receives a suppression by $\epsilon^{2\Delta X^e_{23}}$ or $\epsilon^{2\Delta X^e_{12}}$, respectively. If $X^e_1 = X^e_2 = X^e_3$, all contributions to $(\tilde{\delta}^\ell_{RL})_{12,13,23}$ from the right-handed sleptons disappear, leaving only the left-handed slepton contributions which are suppressed by (m_e/m_μ), (m_e/m_τ) and (m_μ/m_τ), respectively. Since X^ψ_i are quantized, this suppression mechanism does not require any fine tuning of parameters. We stress that this suppression of the flavor violations from A_{IJ}/y_{IJ} relies on the specific feature of the SS SUSY breaking that A_{IJ}/y_{IJ} at the compactification scale are quantized as in (37) if the kink masses are quantized. Expressions for $(\tilde{\delta}^\ell_{RL})_{21,31,32}$ can be obtained from $(\tilde{\delta}^\ell_{RL})_{12,13,23}$ by exchanging $e \leftrightarrow \ell$ for $\Delta X^{e,\ell}$ and $\kappa_{e,\ell}$ together with the exchange $1 \leftrightarrow 2$, $1 \leftrightarrow 3$ and $2 \leftrightarrow 3$, respectively for the first index i in $(\kappa_{e,\ell})_{ij}$, and also moving the phase factor in the first line of the rectangular parenthesis to the second line. A parallel discussion for $(\tilde{\delta}^d_{RL})_{ij}$ and $(\tilde{\delta}^u_{RL})_{ij}$ can be easily understood, so we do not repeat it here.

In the SCKM basis, the right-handed down-type squark masses are given by

$$
(V_d m^{2(\tilde{d})} V^\dagger_d)_{i\bar{j}} \simeq M^2_0\, e^{i(\phi^d_i - \phi^d_j)}\, (\overline{V}_d)_{ik} |N^d_k|^2 \epsilon^{2|N^d_k|} (\overline{V}^\dagger_d)_{kj}.
\tag{65}
$$

The corresponding $\tilde{\delta}$ parameters at the compactification scale can be defined as

$$
(\tilde{\delta}^d_{RR})_{ij} \equiv \frac{(V_d m^{2(\tilde{d})} V^\dagger_d)_{ij}}{|M_{1/2}|^2}.
\tag{66}
$$

We then find

$$
\begin{aligned}
(\tilde{\delta}^d_{RR})_{12} &\simeq -10 e^{i(\phi^d_1-\phi^d_2)}\epsilon^{\Delta X^d_{12}}(\kappa_d)_{11}(\kappa_d)^*_{21} \\
&\quad\times\left[\left(|N^d_2|^2\epsilon^{2|N^d_2|}-|N^d_1|^2\epsilon^{2|N^d_1|}\right)-\frac{(\kappa_d)_{13}(\kappa_d)^*_{23}}{(\kappa_d)_{11}(\kappa_d)^*_{21}}\epsilon^{2\Delta X^d_{23}}\left(|N^d_3|^2\epsilon^{2|N^d_3|}-|N^d_2|^2\epsilon^{2|N^d_2|}\right)\right], \\
(\tilde{\delta}^d_{RR})_{13} &\simeq 10 e^{i(\phi^d_1-\phi^d_3)}\epsilon^{\Delta X^d_{13}}(\kappa_d)_{13}(\kappa_d)^*_{33} \\
&\quad\times\left[\left(|N^d_3|^2\epsilon^{2|N^d_3|}-|N^d_1|^2\epsilon^{2|N^d_1|}\right)+\frac{(\kappa_d)_{12}(\kappa_d)^*_{32}}{(\kappa_d)_{13}(\kappa_d)^*_{33}}\left(|N^d_2|^2\epsilon^{2|N^d_2|}-|N^d_1|^2\epsilon^{2|N^d_1|}\right)\right], \\
(\tilde{\delta}^d_{RR})_{23} &\simeq 10 e^{i(\phi^d_2-\phi^d_3)}\epsilon^{\Delta X^d_{23}}(\kappa_d)_{12}(\kappa_d)^*_{22} \\
&\quad\times\left[\left(|N^d_3|^2\epsilon^{2|N^d_3|}-|N^d_2|^2\epsilon^{2|N^d_2|}\right)-\frac{(\kappa_d)_{21}(\kappa_d)^*_{31}}{(\kappa_d)_{23}(\kappa_d)^*_{33}}\epsilon^{2\Delta X^d_{12}}\left(|N^d_2|^2\epsilon^{2|N^d_2|}-|N^d_1|^2\epsilon^{2|N^d_1|}\right)\right], \qquad (67)
\end{aligned}
$$

where the factor $|N^d_i|^2\epsilon^{2|N^d_i|}$ should be replaced by $1/[2\ln(1/\epsilon)]^2$ for $N^d_i=0$. A suppression of the flavor violations from $m^{2(\tilde{d})}_{i\bar{j}}$ is possible if any pair of $|N^d_i|$ have a common value. In particular, all flavor mixings disappear if $|N^d_1| = |N^d_2| = |N^d_3|$, although this is not favored by the observed quark masses and CKM mixing angles as was noted in the discussion leading to (60). Again this suppression of the flavor violations from $m^2_{I\bar{J}}$ relies on the specific feature of the SS SUSY breaking that $m^2_{I\bar{J}}$ at the compactification scale are quantized as in (37) if the kink masses are quantized. Expressions of $(\tilde{\delta}^\ell_{RR})$ and $(\tilde{\delta}^u_{RR})$ can be obtained by appropriately changing the flavor indices. To obtain $(\tilde{\delta}^\ell_{LL})$, $(\tilde{\delta}^d_{LL})$ and $(\tilde{\delta}^u_{LL})$, one has to remove the phase factors in addition to the necessary change of flavor indices.

The soft terms and mixing parameters discussed above are given at the compactification scale, thus a renormalization group (RG) improvement is required. One approach is a full numerical calculation including the effects of flavor mixing in the RG evolution. This approach would be necessary when the leading flavor mixing comes from the loop effects [7, 8, 20, 21, 22]. However in our case, it is not so useful since the model involves many free parameters. We thus use an approximate analytic solution ignoring the RG effects on the flavor off-diagonal part of the soft parameters [23]. This approximation is reasonably good for the order of magnitude analysis of low energy flavor violations. In the same spirit, we use the mass-insertion approximation [24] to calculate flavor-violating observables at the weak scale, rather than using the more accurate mass-eigenstate formalism.

Using for instance the results of [21], we find that the gaugino masses M_a and the flavor-diagonal sfermion masses $m^{2(\tilde{\psi})}_{i\bar{i}}$ at the weak scale are given by

$$
\begin{aligned}
|M_a^2|/|M_{1/2}|^2 &= 0.16 : 0.67 : 8.5 \\
m^{2(\tilde{q})}_{i\bar{i}}/|M_{1/2}|^2 &\simeq 7.2 : 7.2 : 6.0 \\
m^{2(\tilde{u})}_{i\bar{i}}/|M_{1/2}|^2 &\simeq 6.7 : 6.7 : 4.8 \\
m^{2(\tilde{d})}_{i\bar{i}}/|M_{1/2}|^2 &\simeq 6.7 \quad (i=1,2,3) \\
m^{2(\tilde{\ell})}_{i\bar{i}}/|M_{1/2}|^2 &\simeq 0.53 \quad (i=1,2,3) \\
m^{2(\tilde{e})}_{i\bar{i}}/|M_{1/2}|^2 &\simeq 0.15 \quad (i=1,2,3), \qquad (68)
\end{aligned}
$$

where $M_{1/2} = -\omega/R$ is the universal gaugino mass at the compactification scale and $a = 1,2,3$ stand for $U(1)_Y \times SU(2)_L \times SU(3)_c$, respectively. To get these numerical results, we used the top quark Yukawa coupling $y_t \sim 1$ at the weak scale and ignored all other Yukawa couplings. Similarly the Higgs masses at the weak scale are found to be

$$
m^2_{H_1}/|M_{1/2}|^2 \simeq 0.53\,, \quad m^2_{H_2}/|M_{1/2}|^2 \simeq -3.0. \qquad (69)
$$

We also approximate the Higgsino mass parameter μ as

$$
\begin{aligned}
\mu^2 &= -\frac{M_Z^2}{2} - \frac{m^2_{H_1} - m^2_{H_2}\tan^2\beta}{1-\tan^2\beta} \\
&\simeq -\frac{M_Z^2}{2} + \left(3.0 - \frac{3.5}{1-\tan^2\beta}\right)|M_{1/2}|^2 \simeq 3|M_{1/2}|^2. \qquad (70)
\end{aligned}
$$

The slepton δ parameters at the weak scale are defined in the SCKM basis as [24],

$$
(\delta^\ell_{RL})_{ij} \equiv \frac{A^\ell_{ij} v_d - \mu^* m^\ell_i \delta_{ij}\tan\beta}{\sqrt{m^{2(\tilde{e})}_{i\bar{i}} m^{2(\tilde{\ell})}_{j\bar{j}}}}, \quad (\delta^\ell_{RR})_{ij} \equiv \frac{m^{2(\tilde{e})}_{i\bar{j}}}{\sqrt{m^{2(\tilde{e})}_{i\bar{i}} m^{2(\tilde{e})}_{j\bar{j}}}}, \quad (\delta^\ell_{LL})_{ij} \equiv \frac{m^{2(\tilde{\ell})}_{i\bar{j}}}{\sqrt{m^{2(\tilde{\ell})}_{i\bar{i}} m^{2(\tilde{\ell})}_{j\bar{j}}}}. \qquad (71)
$$

According to (68), these weak scale δ parameters are related to the GUT scale $\tilde{\delta}$ parameters as

$$(\delta^\ell_{RL})_{ij} \simeq 3.5\,(\tilde{\delta}^\ell_{RL})_{ij}, \quad (\delta^\ell_{RR})_{ij} \simeq 6.7\,(\tilde{\delta}^\ell_{RR})_{ij}, \quad (\delta^\ell_{LL})_{ij} \simeq 1.9\,(\tilde{\delta}^\ell_{LL})_{ij} \quad (i \neq j), \tag{72}$$

and similar expressions can be obtained also for the squark δ parameters.

It turns out that the most dangerous low energy flavor violations in our SS SUSY breaking scenario are $\mu \to e\gamma$ and ϵ_K. Let us first discuss the $\mu \to e\gamma$ process. Neglecting the electron mass in the final state, the $\mu \to e\gamma$ branching ratio is given by the sum of two branching ratios with opposite chirality,

$$BR(\mu^+ \to e^+\gamma) \;\simeq\; BR(\mu^+_L \to e^+_R\gamma) + BR(\mu^+_R \to e^+_L\gamma). \tag{73}$$

Assuming the sparticle spectrum of (68) and (70), we find

$$\begin{aligned}
\left[\frac{BR(\mu^+_R \to e^+_L\gamma)}{1.2\times 10^{-11}}\right]^{1/2} \simeq\; & \left(\frac{500\mathrm{GeV}}{|M_{1/2}|}\right)^2 \\
& \times \Bigg| -\frac{(\delta^\ell_{RR})_{12}}{3.1\times 10^{-2}/(e^{-i\theta_\mu}\tan\beta + 2.7)} + \frac{(\delta^\ell_{RR})_{13}(\delta^\ell_{RR})_{32}}{5.1\times 10^{-2}/(e^{-i\theta_\mu}\tan\beta + 3.0)} \\
& \quad + \left(\frac{|M_{1/2}|}{500\mathrm{GeV}}\right)\left(\frac{106\mathrm{MeV}}{m_\mu}\right) \\
& \quad \times \left\{ \frac{(\delta^\ell_{RL})_{12}}{4.8\times 10^{-6}} - \frac{(\delta^\ell_{RL})_{11}(\delta^\ell_{LL})_{12} + (\delta^\ell_{RL})_{13}(\delta^\ell_{LL})_{32}}{6.2\times 10^{-6}} \right. \\
& \quad \left. - \frac{(\delta^\ell_{RR})_{12}(\delta^\ell_{RL})_{22} + (\delta^\ell_{RR})_{13}(\delta^\ell_{RL})_{32}}{8.8\times 10^{-6}} + \frac{(\delta^\ell_{RR})_{13}(\delta^\ell_{RL})_{33}(\delta^\ell_{LL})_{32}}{1.1\times 10^{-5}} \right\} \Bigg|,
\end{aligned} \tag{74}$$

$$\begin{aligned}
\left[\frac{BR(\mu^+_L \to e^+_R\gamma)}{1.2\times 10^{-11}}\right]^{1/2} \simeq\; & \left(\frac{500\mathrm{GeV}}{|M_{1/2}|}\right)^2 \\
& \times \Bigg| \frac{(\delta^\ell_{LL})_{12}}{2.6\times 10^{-2}/(e^{i\theta_\mu}\tan\beta + 0.11)} - \frac{(\delta^\ell_{LL})_{13}(\delta^\ell_{LL})_{32}}{7.8\times 10^{-2}/(e^{i\theta_\mu}\tan\beta - 0.42)} \\
& \quad + \left(\frac{|M_{1/2}|}{500\mathrm{GeV}}\right)\left(\frac{106\mathrm{MeV}}{m_\mu}\right) \\
& \quad \times \left\{ \frac{(\delta^\ell_{LR})_{12}}{4.8\times 10^{-6}} - \frac{(\delta^\ell_{LR})_{11}(\delta^\ell_{RR})_{12} + (\delta^\ell_{LR})_{13}(\delta^\ell_{RR})_{32}}{8.8\times 10^{-6}} \right. \\
& \quad \left. - \frac{(\delta^\ell_{LL})_{12}(\delta^\ell_{LR})_{22} + (\delta^\ell_{LL})_{13}(\delta^\ell_{LR})_{32}}{6.2\times 10^{-6}} + \frac{(\delta^\ell_{LL})_{13}(\delta^\ell_{LR})_{33}(\delta^\ell_{RR})_{32}}{1.1\times 10^{-5}} \right\} \Bigg|,
\end{aligned} \tag{75}$$

where $\theta_\mu \equiv \arg(\mu M_{1/2})$ and the branching ratio is divided by the latest upperbound $BR(\mu \to e\gamma) < 1.2\times 10^{-11}$ [25]. Corresponding analytic formulas can be found *e.g.* in [8] and we expanded chargino and neutralino mixings up to $\mathcal{O}(M_{W,Z}/|M_{1/2}|)$. Here we include two insertions of δ for the RR and LL channels, while only a single insertion of δ is included for the RL channel. Note that not only $(2,1)$ mixings but also some combinations of $(2,3)$ and $(3,1)$ mixings are severely constrained by $\mu \to e\gamma$. Similar expressions for $BR(\tau \to e\gamma)$ and $BR(\tau \to \mu\gamma)$ can be obtained by replacing m_μ with m_τ, changing the generation indices in the δ parameters as $2 \leftrightarrow 3$ and $1 \to 2 \to 3 \to 1$, and multiplying $BR(\tau \to e\nu_\tau\bar{\nu}_e) = 0.178$ and $BR(\tau \to \mu\nu_\tau\bar{\nu}_\mu) = 0.174$, respectively.

Eqs.(64) and (72) suggest that if none of $\Delta X^{e,\ell}_{12}$ vanishes, $|(\delta^\ell_{RL})_{12,21}| = \mathcal{O}(10^{-3}\epsilon^{\Delta X^{e,\ell}_{12}})$ for $|M_{1/2}| \sim 500$ GeV, so it is difficult to satisfy the experimental bound $|(\delta^\ell_{RL})_{12,21}| \lesssim \mathcal{O}(10^{-6})$ under the constraint $(\Delta X^e)_{12} + (\Delta X^\ell)_{12} = 3$, 4 or 5 which comes from m_e/m_μ. A simple mechanism to suppress the $\mu \to e\gamma$ rate is to choose some of the quantized lepton kink masses to be degenerate. For instance, if $X^\ell_1 = X^\ell_2 = X^\ell_3$ or $X^e_1 = X^e_2 = X^e_3$, the $\mu \to e\gamma$ bound can be safely satisfied without severe fine-tuning of the involved $\mathcal{O}(1)$ parameters for $|M_{1/2}| = 500$ GeV. Table I summarizes the possible choices of the flavor charge differences which reproduce the correct charged lepton mass spectrum. The resulting $\mu \to e\gamma$ rate expressed in terms of $(\delta^\ell_{RL})_{12,21}$ for $|M_{1/2}| = 500$GeV. For this, we set all $\mathcal{O}(1)$ parameters, i.e. $\kappa_{e,\ell}$, to be unity, so the results of Table I should be interpreted as a kind of order of magnitude estimate. A double check in the table indicates that the model can safely satisfy the $\mu \to e\gamma$ constraint without any fine tuning of parameters, and a single check means one may need a mild tuning of parameters.

Tables II–III represent the lepton flavor violating rates predicted by the models of Table I. We used Eqs.(74) and (75) with $|M_{1/2}| = 500\,$GeV for $\mu \to e\gamma$, and the analogous formula for $\tau \to \mu\gamma$ or $e\gamma$. In this procedure, we set

all the involved $\mathcal{O}(1)$ parameters to unity, and added all contributions constructively. For $(\delta^\ell_{RL})_{ii}$, we include $\tan\beta$ enhanced F-term contribution,

$$(\delta^\ell_{RL})_{ii} \simeq -0.012 e^{-i\theta_\mu} \tan\beta \left(m^\ell_i/1\text{GeV}\right)\left(500\text{GeV}/|M_{1/2}|\right). \tag{76}$$

Taking into account the sensitivity of next generation experiments [26, 27], the $\mu\to e\gamma$ ($\tau\to e\gamma$ or $\mu\gamma$) branching ratio smaller than 10^{-14} (10^{-9}) for $\tan\beta=10$ is not depicted in the Tables. Models indicated by light color lead to a too rapid $\mu\to e\gamma$ even when the $\mathcal{O}(1)$ parameters are assumed to suppressed by a factor of 1/4 for $|M_{1/2}|=500$ GeV. Note that all branching ratios scale as $|M_{1/2}|^{-4}$, thus the numbers in Tables II−III decrease (increase) by a factor of 1/16 (16) when $|M_{1/2}|=1$ TeV (250 GeV). Many of the models in Tables II and III lead to $\ell_i\to\ell_j\gamma$ which can be explored by future experiments for a reasonable range of $\mathcal{O}(1)$ parameters. Some models already start to overlap with the latest bound $BR(\tau\to\mu\gamma)<3.1\times10^{-7}$ [28] with an ambiguity associated with $\mathcal{O}(1)$ parameters. Different choices of the effective flavor charges predict different patterns of lepton flavor violation, thus a combinatoric analysis of different experiments will be useful for distinguishing models discussed here. In particular, determining the chirality pattern of the processes can provide a crucial information on the model [29]. Note that the chirality pattern for $X^\ell_1=X^\ell_2=X^\ell_3$ is opposite to the case with $X^e_1=X^e_2=X^e_3$ which has the same chirality pattern as the lepton flavor violating decays in seesaw models [7, 8, 9].

Let us now examine the quark sector. It is well known that CP violating parameter ϵ_K in K-$\overline{K}$ mixing is quite sensitive to the supersymmetric extension of the standard model. CP is conserved in the SM if there is no third generation. Consequently, the SM contribution to ϵ_K is suppressed by small CKM mixing angles compared to naive dimensional estimation. However, this is not the case for the supersymmetric models, thus the SUSY contribution to ϵ_K can easily become comparable to the SM contribution. The gluino mediated contribution to ϵ_K, normalized by the experimental value [30], can be summarized as

$$\begin{aligned}\left[\frac{\epsilon_K}{2.282\times10^{-3}}\right] \;\simeq\; e^{i\frac{\pi}{4}}\left(\frac{500\text{GeV}}{|M_{1/2}|}\right)^2 \Bigg[&\frac{\text{Im}[(\delta^d_{LL})^2_{12}]}{(1.5\times10^{-2})^2}+\frac{\text{Im}[(\delta^d_{RR})^2_{12}]}{(1.5\times10^{-2})^2}-\frac{\text{Im}[(\delta^d_{RR})_{12}(\delta^d_{LL})_{12}]}{(2.2\times10^{-4})^2}\\ &+\frac{\text{Im}[(\delta^d_{RL})^2_{12}]+\text{Im}[(\delta^d_{LR})^2_{12}]}{(0.63\times10^{-3})^2}-\frac{\text{Im}[(\delta^d_{RL})_{12}(\delta^d_{LR})_{12}]}{(0.49\times10^{-3})^2}\Bigg],\end{aligned} \tag{77}$$

where we assume again the sparticle mass spectrum of (68). Analytic formulas for the corresponding Wilson coefficients can be found in [24]. Here we followed [31] to estimate the QCD corrections and relevant hadronic matrix elements, and Table IV for the involved phenomenological numbers. The typical size of $(\delta^d_{RL})_{12,21}$ in our SS-mediated SUSY breaking models is $\mathcal{O}(10^{-4})$ as can be seen from the expression of $\tilde\delta^d_{RL}$ analogous to (64) and also (68), therefore the RL insertions do not give an observable contribution to ϵ_K. Also the contributions from $\text{Im}[(\delta^d_{LL,RR})^2_{12}]$ are relatively small because the relevant matrix element does not receive QCD enhancement. In our case, accidental cancellation of relevant mass functions reduces these contributions further. Then the most dominant contribution comes from $\text{Im}[(\delta^d_{RR})_{12}(\delta^d_{LL})_{12}]$. Higher order insertions of δ including the third generation are not important because (67) and (68) show $|(\delta^d_{RR})_{13}(\delta^d_{RR})_{32}|\lesssim|(\delta^d_{RR})_{12}|$ and a similar relation for $(\delta^d_{LL})_{12}$.

Upper half of the Table V summarizes the resulting ϵ_K for the effective flavor charges of (60) providing a best fit to the observed CKM matrix and down-type quark masses under the assumption that $\kappa_{q,d}$ and $e^{i(\phi^d_1-\phi^d_2)}$ are complex in general. The results are expressed in terms of the δ parameters normalized by their values saturating the observed $|\epsilon_K|$. Again here we choose $|M_{1/2}|\sim 500$ GeV and assume that the complex parameters $\kappa_{q,d}$ and the phases ϕ^d_i are all of order unity. For these four models, the contribution from $\text{Im}[(\delta^d_{RR})_{12}(\delta^d_{LL})_{12}]$ exceeds the observed value by one or two orders of magnitude, thus one needs a fine tuning of $\mathcal{O}(10^{-1}\sim10^{-2})$ for the involved parameters. Even when $\kappa_{q,d}$ and $e^{i(\phi^d_1-\phi^d_2)}$ are all real, the situation is not improved much. In the presence of the KM phase, the RG evolution from M_c to M_Z generates a CP violating phase in $(\delta^d_{LL})_{12}$:

$$\begin{aligned}\text{Im}[(\delta^d_{LL})_{12}] \;&\sim\; \frac{1}{\sqrt{m^{2(\tilde q)}_{11}m^{2(\tilde q)}_{22}}}\text{Im}[(A^\dagger_u A_u)_{12}]\frac{1}{(4\pi)^2}\ln\left(\frac{M_Z^2}{M_c^2}\right)\\ &\sim\; -6\times10^{-2}\left(\frac{|M_{1/2}|}{v_u}\right)^2\text{Im}[(\tilde\delta^u_{RL})^*_{31}(\tilde\delta^u_{RL})_{32}]\\ &\sim\; -10^{-3}\,\text{Im}[(V_{CKM}\bar V_q)_{13}(V_{CKM}\bar V_q)^*_{23}]/\epsilon^{X^q_1+X^q_2-2X^q_3},\end{aligned} \tag{78}$$

where $v_u=\langle H^0_2\rangle\simeq 174\sin\beta$ GeV. This RG induced contribution is numerically only a factor few smaller than the direct contribution at M_c coming from complex $\kappa_{q,d}$ and $e^{i(\phi^d_1-\phi^d_2)}$.

$\hat{y}^\ell_i/\hat{y}^\ell_3=\mathcal{O}(\epsilon^5,\epsilon,1)$					$\hat{y}^\ell_i/\hat{y}^\ell_3=\mathcal{O}(\epsilon^6,\epsilon,1)$				
ΔX^e_{i3}	ΔX^ℓ_{i3}	$\frac{\lvert(\delta^\ell_{RL})_{12}\rvert}{4.8\times10^{-6}}$	$\frac{\lvert(\delta^\ell_{LR})_{12}\rvert}{4.8\times10^{-6}}$		ΔX^e_{i3}	ΔX^ℓ_{i3}	$\frac{\lvert(\delta^\ell_{RL})_{12}\rvert}{4.8\times10^{-6}}$	$\frac{\lvert(\delta^\ell_{LR})_{12}\rvert}{4.8\times10^{-6}}$	
–	–	–	–		$(6,1,0)$	$(0,0,0)$	0.80	0.040	$\surd\surd$
$(5,1,0)$	$(0,0,0)$	3.2	0.015	$\surd$	$(5,1,0)$	$(1,0,0)$	3.2	100	
$(4,1,0)$	$(1,0,0)$	12	100		$(4,1,0)$	$(2,0,0)$	12	40	
$(3,1,0)$	$(2,0,0)$	40	40		$(3,1,0)$	$(3,0,0)$	40	12	
$(2,1,0)$	$(3,0,0)$	100	12		$(2,1,0)$	$(4,0,0)$	100	3.2	
$(1,1,0)$	$(4,0,0)$	20	3.2		$(1,1,0)$	$(5,0,0)$	20	0.80	
–	–	–	–		$(5,0,0)$	$(1,1,0)$	0.80	20	
$(4,0,0)$	$(1,1,0)$	3.2	20		$(4,0,0)$	$(2,1,0)$	3.2	100	
$(3,0,0)$	$(2,1,0)$	12	100		$(3,0,0)$	$(3,1,0)$	12	40	
$(2,0,0)$	$(3,1,0)$	40	40		$(2,0,0)$	$(4,1,0)$	40	12	
$(1,0,0)$	$(4,1,0)$	100	12		$(1,0,0)$	$(5,1,0)$	100	3.2	
$(0,0,0)$	$(5,1,0)$	0.015	3.2	$\surd$	$(0,0,0)$	$(6,1,0)$	0.040	0.80	$\surd\surd$
$\hat{y}^\ell_i/\hat{y}^\ell_3=\mathcal{O}(\epsilon^5,\epsilon^2,1)$					$\hat{y}^\ell_i/\hat{y}^\ell_3=\mathcal{O}(\epsilon^6,\epsilon^2,1)$				
ΔX^e_{i3}	ΔX^ℓ_{i3}	$\frac{\lvert(\delta^\ell_{RL})_{12}\rvert}{4.8\times10^{-6}}$	$\frac{\lvert(\delta^l_{LR})_{12}\rvert}{4.8\times10^{-6}}$		ΔX^e_{i3}	ΔX^ℓ_{i3}	$\frac{\lvert(\delta^l_{RL})_{12}\rvert}{4.8\times10^{-6}}$	$\frac{\lvert(\delta^\ell_{LR})_{12}\rvert}{4.8\times10^{-6}}$	
–	–	–	–		$(6,2,0)$	$(0,0,0)$	3.2	0.60	$\surd$
$(5,2,0)$	$(0,0,0)$	12	0.057		$(5,2,0)$	$(1,0,0)$	12	100	
$(4,2,0)$	$(1,0,0)$	40	100		$(4,2,0)$	$(2,0,0)$	40	40	
$(3,2,0)$	$(2,0,0)$	100	40		$(3,2,0)$	$(3,0,0)$	100	12	
$(2,2,0)$	$(3,0,0)$	1.6	12		$(2,2,0)$	$(4,0,0)$	1.6	3.2	$\surd$
–	–	–	–		$(5,1,0)$	$(1,1,0)$	3.2	20	
$(4,1,0)$	$(1,1,0)$	12	20		$(4,1,0)$	$(2,1,0)$	12	100	
$(3,1,0)$	$(2,1,0)$	40	100		$(3,1,0)$	$(3,1,0)$	40	40	
$(2,1,0)$	$(3,1,0)$	100	40		$(2,1,0)$	$(4,1,0)$	100	12	
$(1,1,0)$	$(4,1,0)$	20	12		$(1,1,0)$	$(5,1,0)$	20	3.2	
–	–	–	–		$(4,0,0)$	$(2,2,0)$	3.2	1.6	$\surd$
$(3,0,0)$	$(2,2,0)$	12	1.6		$(3,0,0)$	$(3,2,0)$	12	100	
$(2,0,0)$	$(3,2,0)$	40	100		$(2,0,0)$	$(4,2,0)$	40	40	
$(1,0,0)$	$(4,2,0)$	100	40		$(1,0,0)$	$(5,2,0)$	100	12	
$(0,0,0)$	$(5,2,0)$	0.057	12		$(0,0,0)$	$(6,2,0)$	0.60	3.2	$\surd$

TABLE I: Lepton mass hierarchy vs constraint from $\mu\to e\gamma$. Here $(\delta^\ell_{RL})_{12,21}$ is divided by the values saturating $BR(\mu\to e\gamma)=1.2\times10^{-11}$ for $|M_{1/2}|=500$GeV. A double check indicates that the model can safely satisfy the $\mu\to e\gamma$ constraint and a single check means one may need a mild tuning of $\mathcal{O}(1)$ parameters.

If we relax the condition (60) for the best fit to the quark masses and CKM matrix, we can choose $N^d_1=N^d_2=N^d_3$ for the down-type quark singlets, which would make all $(\delta^d_{RR})_{ij}$ disappear. Note that the SS-induced $m^{2(\tilde{d})}_{i\bar{j}}$ at the compactification scale are universal if N^d_i are all degenerate. Lower half of the Table V shows that in this case there is no observable deviation of ϵ_K. However in this case, in order to produce the correct quark masses and CKM matrix, we have to assume that some boundary Yukawa couplings are abnormally large (or small) by a factor of $4\sim5$ $(0.2\sim0.3)$ compared to the values suggested by the naive dimensional analysis [15]. For instance, the model with $X^d_i=(3,2,0)$ requires that the boundary Yukawa coupling $\tilde{\lambda}_{11}$ is smaller than the naively expected value by a factor of 0.2.

For the sparticle spectrum (68), the gluino contribution to the K^0-$\overline{K}^0$ and $B^0_{d,s}$-$\overline{B}^0_{d,s}$ mass differences in SS SUSY breaking scenario are given by

$$\left[\frac{\Delta M_K}{5.30\text{ns}^{-1}\hbar}\right]\simeq\left(\frac{500\text{GeV}}{|M_{1/2}|}\right)^2\left[\frac{\text{Re}[(\delta^d_{LL})^2_{12}]}{(1.8\times10^{-1})^2}+\frac{\text{Re}[(\delta^d_{RR})^2_{12}]}{(1.9\times10^{-1})^2}-\frac{\text{Re}[(\delta^d_{RR})_{12}(\delta^d_{LL})_{12}]}{(2.8\times10^{-3})^2}\right.$$
$$\left.+\frac{\text{Re}[(\delta^d_{RL})^2_{12}]+\text{Re}[(\delta^d_{LR})^2_{12}]}{(0.78\times10^{-2})^2}-\frac{\text{Re}[(\delta^d_{RL})_{12}(\delta^d_{LR})_{12}]}{(0.60\times10^{-2})^2}\right],$$

$y_i^\ell = \mathcal{O}(\epsilon^7, \epsilon^3, \epsilon^2)$				
N_i^e	N_i^ℓ	$BR(\mu_R^+ \to e_L^+, \gamma)$	$BR(\tau_R^+ \to e_L^+, \gamma)$	$BR(\tau_R^+ \to \mu_L^+, \gamma)$
$(5,1,-2)$	$(2,2,2)$	$1.5(1+0.14t_\beta)^2 \times 10^{-10}$	$-$	$2.4(1+0.091t_\beta)^2 \times 10^{-8}$
$(5,1,-1)$	$(2,2,2)$	$1.6(1+0.14t_\beta)^2 \times 10^{-10}$	$-$	2.1×10^{-8}
$(6,2,1)$	$(1,1,1)$	$1.3(1+0.026t_\beta)^2 \times 10^{-10}$	$-$	$4.4(1+0.42t_\beta)^2 \times 10^{-8}$
$(7,3,2)$	$(-1,-1,-1)$	1.2×10^{-10}	$-$	$2.4(1+0.091t_\beta)^2 \times 10^{-8}$
$(-1,-1,-1)$	$(7,3,2)$	$-$	$-$	$-$
$(1,1,1)$	$(6,2,1)$	$-$	$-$	$-$
$(2,2,2)$	$(5,1,-1)$	$-$	$-$	$-$
$(2,2,2)$	$(5,1,-2)$	$-$	$-$	$-$
N_i^e	N_i^ℓ	$BR(\mu_L^+ \to e_R^+, \gamma)$	$BR(\tau_L^+ \to e_R^+, \gamma)$	$BR(\tau_L^+ \to \mu_R^+, \gamma)$
$(5,1,-2)$	$(2,2,2)$	$-$	$-$	$-$
$(5,1,-1)$	$(2,2,2)$	$-$	$-$	$-$
$(6,2,1)$	$(1,1,1)$	$-$	$-$	$4.6t_\beta^2 \times 10^{-11}$
$(7,3,2)$	$(-1,-1,-1)$	$-$	$-$	$-$
$(-1,-1,-1)$	$(7,3,2)$	1.2×10^{-10}	$-$	$2.1(1+0.060t_\beta)^2 \times 10^{-8}$
$(1,1,1)$	$(6,2,1)$	$1.2(1+0.015t_\beta)^2 \times 10^{-10}$	$-$	$2.1(1+0.37t_\beta)^2 \times 10^{-8}$
$(2,2,2)$	$(5,1,-1)$	$1.3(1+0.092t_\beta)^2 \times 10^{-10}$	$-$	2.1×10^{-8}
$(2,2,2)$	$(5,1,-2)$	$1.2(1+0.093t_\beta)^2 \times 10^{-10}$	$-$	$2.1(1+0.015t_\beta)^2 \times 10^{-8}$
$y_i^\ell = \mathcal{O}(\epsilon^8, \epsilon^3, \epsilon^2)$				
N_i^e	N_i^ℓ	$BR(\mu_R^+ \to e_L^+, \gamma)$	$BR(\tau_R^+ \to e_L^+, \gamma)$	$BR(\tau_R^+ \to \mu_L^+, \gamma)$
$(6,1,-2)$	$(2,2,2)$	$9.2(1+0.11t_\beta)^2 \times 10^{-12}$	$-$	$2.4(1+0.091t_\beta)^2 \times 10^{-8}$
$(6,1,-1)$	$(2,2,2)$	$9.4(1+0.11t_\beta)^2 \times 10^{-12}$	$-$	2.1×10^{-7}
$(7,2,1)$	$(1,1,1)$	$8.2(1+0.021t_\beta)^2 \times 10^{-12}$	$-$	$4.4(1+0.42t_\beta)^2 \times 10^{-8}$
$(8,3,2)$	$(-1,-1,-1)$	7.8×10^{-12}	$-$	$2.4(1+0.091t_\beta)^2 \times 10^{-8}$
$(-1,-1,-1)$	$(8,3,2)$	$-$	$-$	$-$
$(1,1,1)$	$(7,2,1)$	$-$	$-$	$-$
$(2,2,2)$	$(6,1,-1)$	$-$	$-$	$-$
$(2,2,2)$	$(6,1,-2)$	$-$	$-$	$-$
N_i^e	N_i^ℓ	$BR(\mu_L^+ \to e_R^+, \gamma)$	$BR(\tau_L^+ \to e_R^+, \gamma)$	$BR(\tau_L^+ \to \mu_R^+, \gamma)$
$(6,1,-2)$	$(2,2,2)$	$-$	$-$	$-$
$(6,1,-1)$	$(2,2,2)$	$-$	$-$	$-$
$(7,2,1)$	$(1,1,1)$	$-$	$-$	$4.6t_\beta^2 \times 10^{-11}$
$(8,3,2)$	$(-1,-1,-1)$	$-$	$-$	$-$
$(-1,-1,-1)$	$(8,3,2)$	7.7×10^{-12}	$-$	$2.1(1+0.060t_\beta)^2 \times 10^{-8}$
$(1,1,1)$	$(7,2,1)$	$7.8(1+0.012t_\beta)^2 \times 10^{-12}$	$-$	$2.1(1+0.37t_\beta)^2 \times 10^{-8}$
$(2,2,2)$	$(6,1,-1)$	$7.8(1+0.074t_\beta)^2 \times 10^{-12}$	$-$	2.1×10^{-8}
$(2,2,2)$	$(6,1,-2)$	$7.7(1+0.075t_\beta)^2 \times 10^{-12}$	$-$	$2.1(1+0.060t_\beta)^2 \times 10^{-8}$

TABLE II: Predictions of lepton flavor violating rates for $|M_{1/2}| = 500\text{GeV}$. Here all parameters of $\mathcal{O}(1)$ are set to 1 and all leading contributions are added constructively, so the actual rates can be somewhat smaller than the numbers in the table. The $\mu \to e\gamma$ ($\tau \to e\gamma$ or $\mu\gamma$) branching ratio smaller than 10^{-14} (10^{-9}) for $t_\beta \equiv \tan\beta = 10$ is omitted. Note that the branching ratios scale as $|M_{1/2}|^{-4}$.

$$\begin{aligned}
\left[\frac{\Delta M_{B_d}}{0.489\text{ps}^{-1}\hbar}\right] &\simeq \left(\frac{500\text{GeV}}{|M_{1/2}|}\right)^2 \left| \frac{(\delta^d_{LL})^2_{13}}{(3.5\times 10^{-1})^2} + \frac{(\delta^d_{RR})^2_{13}}{(3.4\times 10^{-1})^2} - \frac{(\delta^d_{RR})_{13}(\delta^d_{LL})_{13}}{(3.2\times 10^{-2})^2} \right. \\
&\qquad \left. + \frac{(\delta^d_{RL})^2_{13}}{(0.65\times 10^{-1})^2} + \frac{(\delta^d_{LR})^2_{13}}{(0.66\times 10^{-1})^2} - \frac{(\delta^d_{RL})_{13}(\delta^d_{LR})_{13}}{(0.71\times 10^{-1})^2} \right|, \\
\left[\frac{\Delta M_{B_s}}{13.1\text{ps}^{-1}\hbar}\right] &\simeq \left(\frac{500\text{GeV}}{|M_{1/2}|}\right)^2 \left| \frac{(\delta^d_{LL})^2_{23} + (\delta^d_{RR})^2_{23}}{(1.5)^2} - \frac{(\delta^d_{RR})_{23}(\delta^d_{LL})_{23}}{(1.4\times 10^{-1})^2} \right.
\end{aligned}$$

$y^\ell_i = \mathcal{O}(\epsilon^8, \epsilon^4, \epsilon^2)$				
N^e_i	N^ℓ_i	$BR(\mu^+_R \to e^+_L, \gamma)$	$BR(\tau^+_R \to e^+_L, \gamma)$	$BR(\tau^+_R \to \mu^+_L, \gamma)$
$(6,2,-2)$	$(2,2,2)$	$1.3(1+0.024t_\beta)^2 \times 10^{-10}$	–	3.3×10^{-9}
$(6,2,-1)$	$(2,2,2)$	$1.3(1+0.024t_\beta)^2 \times 10^{-10}$	–	$5.0(1+0.25t_\beta)^2 \times 10^{-9}$
$(7,3,1)$	$(1,1,1)$	1.2×10^{-10}	–	$5.0(1+0.25t_\beta)^2 \times 10^{-9}$
$(8,4,2)$	$(-1,-1,-1)$	1.2×10^{-10}	–	$2.6(1+0.047t_\beta)^2 \times 10^{-9}$
$(2,2,-2)$	$(6,2,2)$	3.1×10^{-11}	3.4×10^{-9}	3.3×10^{-9}
$(2,2,-1)$	$(6,2,2)$	$8.4(1+0.27t_\beta)^2 \times 10^{-11}$	$5.2(1+0.25t_\beta)^2 \times 10^{-9}$	$5.1(1+0.25t_\beta)^2 \times 10^{-9}$
$(3,3,1)$	$(5,1,1)$	$8.4(1+0.27t_\beta)^2 \times 10^{-11}$	$5.2(1+0.25t_\beta)^2 \times 10^{-9}$	$5.1(1+0.25t_\beta)^2 \times 10^{-9}$
$(4,4,2)$	$(4,-1,-1)$	$3.4(1+0.051t_\beta)^2 \times 10^{-11}$	$3.7(1+0.047t_\beta)^2 \times 10^{-9}$	$3.6(1+0.047t_\beta)^2 \times 10^{-9}$
$(4,4,2)$	$(4,-2,-2)$	$3.4(1+0.053t_\beta)^2 \times 10^{-11}$	$3.7(1+0.047t_\beta)^2 \times 10^{-9}$	$3.6(1+0.047t_\beta)^2 \times 10^{-9}$
$(4,-2,-2)$	$(4,4,2)$	$1.9(1+0.12t_\beta)^2 \times 10^{-10}$	–	–
$(4,-1,-1)$	$(4,4,2)$	$3.4(1+0.10t_\beta)^2 \times 10^{-10}$	–	–
$(5,1,1)$	$(3,3,1)$	$4.9(1+0.12t_\beta)^2 \times 10^{-10}$	–	–
$(6,2,2)$	$(2,2,-1)$	$2.7(1+0.026t_\beta)^2 \times 10^{-10}$	–	–
$(6,2,2)$	$(2,2,-2)$	$1.5(1+0.022t_\beta)^2 \times 10^{-10}$	–	–
$(-1,-1,-1)$	$(8,4,2)$	–	–	–
$(1,1,1)$	$(7,3,1)$	–	–	–
$(2,2,2)$	$(6,2,-1)$	–	–	–
$(2,2,2)$	$(6,2,-2)$	–	–	–
N^e_i	N^ℓ_i	$BR(\mu^+_L \to e^+_R, \gamma)$	$BR(\tau^+_L \to e^+_R, \gamma)$	$BR(\tau^+_L \to \mu^+_R, \gamma)$
$(6,2,-1)$	$(2,2,2)$	–	–	–
$(6,2,-1)$	$(2,2,2)$	–	–	–
$(7,3,1)$	$(1,1,1)$	–	–	–
$(8,4,2)$	$(-1,-1,-1)$	–	–	–
$(2,2,-2)$	$(6,2,2)$	$1.3(1+0.014t_\beta)^2 \times 10^{-10}$	–	–
$(2,2,-1)$	$(6,2,2)$	$5.0(1+0.015t_\beta)^2 \times 10^{-10}$	–	–
$(3,3,1)$	$(5,1,1)$	$5.8(1+0.082t_\beta)^2 \times 10^{-10}$	–	–
$(4,4,2)$	$(4,-1,-1)$	$2.3(1+0.079t_\beta)^2 \times 10^{-10}$	–	–
$(4,4,2)$	$(4,-2,-2)$	$1.8(1+0.32t_\beta)^2 \times 10^{-10}$	–	–
$(4,-2,-2)$	$(4,4,2)$	$3.1(1+0.032t_\beta)^2 \times 10^{-11}$	$3.4(1+0.030t_\beta)^2 \times 10^{-9}$	$3.3(1+0.030t_\beta)^2 \times 10^{-9}$
$(4,-1,-1)$	$(4,4,2)$	$3.1(1+0.032t_\beta)^2 \times 10^{-11}$	$3.4(1+0.030t_\beta)^2 \times 10^{-9}$	$3.3(1+0.030t_\beta)^2 \times 10^{-9}$
$(5,1,1)$	$(3,3,1)$	$3.1(1+0.20t_\beta)^2 \times 10^{-11}$	$3.4(1+0.19t_\beta)^2 \times 10^{-9}$	$3.3(1+0.19t_\beta)^2 \times 10^{-9}$
$(6,2,2)$	$(2,2,-1)$	$3.1(1+0.20t_\beta)^2 \times 10^{-11}$	$3.4(1+0.19t_\beta)^2 \times 10^{-9}$	$3.3(1+0.19t_\beta)^2 \times 10^{-9}$
$(6,2,2)$	$(2,2,-2)$	3.1×10^{-11}	3.4×10^{-9}	3.3×10^{-9}
$(-1,-1,-1)$	$(8,4,2)$	1.2×10^{-10}	–	$3.3(1+0.030t_\beta)^2 \times 10^{-9}$
$(1,1,1)$	$(7,3,1)$	1.2×10^{-10}	–	$3.3(1+0.19t_\beta)^2 \times 10^{-9}$
$(2,2,2)$	$(6,2,-1)$	$1.2(1+0.015t_\beta)^2 \times 10^{-10}$	–	$3.3(1+0.19t_\beta)^2 \times 10^{-9}$
$(2,2,2)$	$(6,2,-2)$	$1.2(1+0.015t_\beta)^2 \times 10^{-10}$	–	3.3×10^{-9}

TABLE III: Predictions of lepton flavor violating rates. Models indicated by light color can not satisfy the current experimental bound on $\mu \to e\gamma$ even when the involved $\mathcal{O}(1)$ parameters are assumed to be suppressed by a factor of 1/4 or $|M_{1/2}| = 1$ TeV.

$$+\frac{(\delta^d_{RL})^2_{23} + (\delta^d_{LR})^2_{23}}{(2.8\times 10^{-1})^2} - \frac{(\delta^d_{RL})_{23}(\delta^d_{LR})_{23}}{(3.0\times 10^{-1})^2}\Bigg| . \quad (79)$$

Here we followed Ref.[24] for the involved Wilson coefficients, Ref.[31] for the QCD corrections and hadronic matrix elements, and Table IV for the involved phenomenological numbers. The results are then well below the experimental values ($\Delta M_{K,B_d}$) or the latest upper bound (ΔM_{B_s}) [30], typically less than 1 %. In fact, for ΔM_{B_d}, there can be few % SUSY contributions which were not included in (79) as they come from the RG effects involving the top quark Yukawa coupling similarly to the effects of (78). Flavor violating soft parameters can affect also ϵ'/ϵ_K even when they do not contain any new CP violating phase. We have estimated the gluino contribution to ϵ'/ϵ_K in SS SUSY

$\alpha_s(M_Z)$	M_K	ΔM_K	F_K	M_{B_d}	F_{B_d}	M_{B_s}	F_{B_s}/F_{B_d}
0.119	494 MeV	$5.30 ns^{-1}\hbar$	160 MeV	5.28 GeV	195 MeV	5.37 GeV	1.21

TABLE IV: Input parameters used in the mass-insertion formulas. For the hadronic matrix elememts, we follow Ref. [31].

N_i^q	N_i^d	$\frac{\mathrm{Im}[(\delta^d_{LL})^2_{12}]}{(1.5\times10^{-2})^2}$	$\frac{\mathrm{Im}[(\delta^d_{RR})^2_{12}]}{(1.5\times10^{-2})^2}$	$\frac{\mathrm{Im}[(\delta^d_{RR})_{12}(\delta^d_{LL})_{12}]}{(2.2\times10^{-4})^2}$
$(3,2,-1)$	$(3,2,2)$	1.5×10^{-2}	7.1×10^{-1}	466
$(3,2,-1)$	$(4,2,2)$	1.5×10^{-2}	2.8×10^{-2}	93
$(3,2,-1)$	$(3,3,2)$	1.5×10^{-2}	2.8×10^{-2}	93
$(3,2,-1)$	$(4,3,2)$	1.5×10^{-2}	5.7×10^{-3}	42
$(3,2,-1)$	$(2,2,2)$	1.5×10^{-2}	0	0
$(4,3,-1)$	$(2,2,2)$	1.2×10^{-4}	0	0
$(4,2,-1)$	$(2,2,2)$	6.0×10^{-4}	0	0

TABLE V: Quark mass hierarchy vs ϵ_K. Here δ^d 's are divided by the values saturating $\epsilon_K = 2.282\times10^{-3}$ for $|M_{1/2}| = 500$GeV.

breaking scenario, and find that it is at most comparable to the SM contribution for a reasonable range of the involved parameters. Because the consensus on the SM contribution to ϵ'/ϵ_K has not been achieved yet [32], we can not derive any meaningful constraint from this result.

Let us finally consider the SUSY contribution to $b\to s\gamma$ in our models. The branching ratio can be approximated by

$$BR[\overline{B}\to X_s,\gamma]^{\rm SUSY} \simeq \Delta BR(b_R\to s_L,\gamma)^{int} + BR(b_L\to s_R,\gamma), \tag{80}$$

where the first term denotes interference with the SM contribution and the second term comes from the operators of opposite chirality to the SM ones. The gluino contributions normalized by the latest world average [33] are given by

$$\begin{aligned}\left[\frac{\Delta BR(b_R\to s_L,\gamma)^{int}}{3.34\,(\pm0.38)\times10^{-4}}\right] &\simeq \left(\frac{500\mathrm{GeV}}{|M_{1/2}|}\right)\left\{\frac{\mathrm{Re}[(\delta^d_{LR})_{23}]}{0.021} - \frac{\mathrm{Re}[(\delta^d_{LL})_{21}(\delta^d_{LR})_{13}]}{0.040} - \frac{\mathrm{Re}[(\delta^d_{LL})_{23}(\delta^d_{LR})_{33}]}{0.043}\right.\\ &\qquad\left. - \frac{\mathrm{Re}[(\delta^d_{LR})_{21}(\delta^d_{RR})_{13}] + \mathrm{Re}[(\delta^d_{LR})_{22}(\delta^d_{RR})_{23}]}{0.041}\right\}, \end{aligned}\tag{81}$$

$$\begin{aligned}\left[\frac{BR(b_L\to s_R,\gamma)}{3.34\,(\pm0.38)\times10^{-4}}\right]^{1/2} &\simeq \left(\frac{500\mathrm{GeV}}{|M_{1/2}|}\right)\left|-\frac{(\delta^d_{RL})_{23}}{0.042} + \frac{(\delta^d_{RR})_{21}(\delta^d_{RL})_{13} + (\delta^d_{RR})_{23}(\delta^d_{RL})_{33}}{0.082}\right.\\ &\qquad\left. + \frac{(\delta^d_{RL})_{21}(\delta^d_{LL})_{13} + (\delta^d_{RL})_{22}(\delta^d_{LL})_{23}}{0.080}\right|, \end{aligned}\tag{82}$$

for the sparticle spectrum (68). Here we followed Ref.[34] to estimate the branching ratio from the relevant Wilson coefficients at m_t and included the leading–order gluino contribution to $C_{4,7,8}$ and their chirality partners. The Ref.[34] quotes the SM contribution as $BR[\overline{B}\to X_s\gamma]_{E_0>\frac{m_b}{20}} = 3.70\pm0.30\times10^{-4}$. This overlaps with the experimental value within the 1 σ errors $\sim 10\%$. In our models, $|\delta^d_{RL,LR}|$ is at most 10^{-3}. Therefore their contributions to the interference term is only a few % level. A potentially dangerous contribution may come from $(\delta^d_{LL})_{23}(\delta^d_{LR})_{33}$ associated with the F-term contribution to $(\delta^d_{LR})_{33}$ which can be estimated as

$$(\delta^d_{LR})_{33} \simeq -1.5\times10^{-3} e^{i\theta_\mu}\tan\beta(500\mathrm{GeV}/|M_{1/2}|)$$

for a moderate value $\tan\beta$. For the models listed in Table V, (δ^d_{LL}) is at most a few % for $|M_{1/2}|\sim 500$ GeV, and then its contribution to $b\to s\gamma$ is also below a few %. The pure SUSY contribution from the opposite chirality is even more negligible because it scales quadratically with the SUSY amplitude, so typically gives a correction of $\mathcal{O}(10^{-3})$. Consequently the gluino mediated contribution to $b\to s\gamma$ does not give any meaningful constraint for our models with the current experimental and theoretical accuracy.

IV. CONCLUSION

Quasi-localization of matter fields in extra dimension is an elegant mechanism to generate hierarchical 4D Yukawa couplings. Extra dimension provides also an attractive way to break SUSY by boundary conditions as originally proposed by Scherk and Schwarz. In this paper, we have examined some physical consequences of implementing the quasi-localization of matter zero modes and the SS SUSY breaking simultaneously within 5D orbifold field theories. In this case, the radion corresponds to a flavon to generate the flavor hierarchy and at the same time plays the role of the messenger of supersymmetry breaking. As a consequence, the resulting soft scalar masses and trilinear A-parameters of matter zero modes at the compactification scale are highly flavor-dependent, thereby can lead to dangerous flavor violations at low energy scales. The shape of soft parameters implies also that the compactification scale should be much higher than the weak scale in order for the model to be phenomenologically viable.

We have computed the soft parameters of quasi-localized matter fields induced by the SS boundary condition in generic 5D orbifold SUGRA. It is shown explicitly that the zero mode soft parameters from the SS boundary condition are same as the ones induced by the radion F-component in 4D effective SUGRA, and thus our analysis applies to any SUSY breaking mechanism giving a sizable F-component of the radion superfield, e.g. the hidden gaugino condensation model.

In 5D orbifold SUGRA, quasi-localization of matter zero modes are governed by the kink masses of matter hypermultiplets, $M_I\epsilon(y)$, which have quantized-values if the graviphoton and/or $U(1)_{FI}$ gauge charges are quantized. An important feature of the SS SUSY breaking or the radion-mediated SUSY breaking is that, if the kink masses are quantized, the resulting soft scalar masses and the trilinear scalar couplngs (divided by the corresponding Yukawa couplings) at the compactification scale are quantized also in the leading approximation. This feature provides a natural mechanism to suppress dangerous flavor violations since the flavor violating amplitudes appear in a form $f(M_I) - f(M_J)$, thus are canceled when some of the quantized kink masses are degenerate.

We analyzed in detail the low energy flavor violations in SS-dominated supersymmetry breaking scenario under the assumption that the compactification scale M_c is close to the grand unification scale $\sim 2 \times 10^{16}$ GeV and the 4D effective theory below M_c is the minimal supersymmetric standard model. We find that many of the low energy flavor violations are appropriately suppressed, however generically ϵ_K and $\mu \to e\gamma$ can be dangerous if the SS boundary condition is the major source of SUSY breaking. Assuming that the hypermultiplet kink masses are quantized, the $\mu \to e\gamma$ bound can be satisfied for a reasonable range of the involved continuous parameters if either the $SU(2)_L$ doublet lepton kink masses or the $SU(2)_L$ singlet lepton kink masses are flavor-independent. These two possibilities are clearly distinguished by the predicted chirality pattern of the lepton flavor violating decays. The chirality structure for degenerate $SU(2)_L$-doublet lepton kink masses is opposite to the other case with degenerate $SU(2)_L$-singlet lepton kink masses which has the same chirality structure as the lepton flavor violating decays induced by the right-handed neutrino Yukawa couplings in seesaw models [7, 8, 9]. The SUSY contribution to ϵ_K can be similarly suppressed by choosing the quantized kink masses of down-type quarks to be degenerate. However in this case, to get the correct quark mass spectrum and CKM mixing angles, one needs to assume that some boundary Yukawa couplings are abnormally large (or small) by a factor of $4 \sim 5$ ($0.2 \sim 0.3$) compared to the values suggested by the naive dimensional analysis.

Acknowledgements

This work is supported by KRF PBRG 2002-070-C00022.

[1] H. P. Nilles, Phys. Rept. **110**, 1 (1984); H. E. Haber and G. L. Kane, Phys. Rept. **117**, 75 (1985).

[2] J. Scherk and J. H. Schwarz, Phys. Lett. B **82**, 60 (1979).

[3] D. E. Kaplan and N. Weiner, hep-ph/0108001; G. von Gersdorff and M. Quiros, Phys. Rev. D **65**, 064016 (2002) [hep-th/0110132]; G. von Gersdorff, M. Quiros and A. Riotto, Nucl. Phys. B **634**, 90 (2002) [hep-th/0204041]; A. Delgado, G. von Gersdorff and M. Quiros, JHEP **0212**, 002 (2002) [hep-th/0210181]; M. Quiros, hep-ph/0302189.

[4] N. Arkani-Hamed and M. Schmaltz, Phys. Rev. D **61**, 033005 (2000) [hep-ph/9903417]; E. A. Mirabelli and M. Schmaltz, Phys. Rev. D **61**, 113011 (2000) [hep-ph/9912265]; D. E. Kaplan and T. M. P. Tait, JHEP **0111**, 051 (2001) [hep-ph/0110126]; M. Kakizaki and M. Yamaguchi, Int. J. Mod. Phys. A **19**, 1715 (2004) [hep-ph/0110266]; N. Haba and N. Maru, Phys. Rev. D **66**, 055005 (2002) [hep-ph/0204069]; A. Hebecker and J. March-Russell, Phys. Lett. B **541**, 338 (2002) [hep-ph/0205143]; K. Choi, D. Y. Kim, I. W. Kim and T. Kobayashi, Eur. Phys. J **C35**, 267 (2004) [hep-ph/0305024]; G. von Gersdorff, L. Pilo, M. Quiros, D. A. J. Rayner and A. Riotto, Phys. Lett. B **580**, 93 (2004) [hep-ph/0305218]; K. Choi, I. W. Kim and W. Y. Song, Nucl. Phys. B **687**, 101 (2004) [hep-ph/0307365]; G. von Gersdorff, M. Quiros and A. Riotto, hep-th/0310190;

[5] R. Barbieri, R. Contino, P. Creminelli, R. Rattazzi and C. A. Scrucca, Phys. Rev. D **66**, 024025 (2002) [hep-th/0203039]; S. Groot Nibbelink, H. P. Nilles and M. Olechowski, Phys. Lett. B **536**, 270 (2002) [hep-th/0203055]; S. Groot Nibbelink, H. P. Nilles and M. Olechowski, Nucl. Phys. B **640**, 171 (2002) [hep-th/0205012]; D. Marti and A. Pomarol, Phys. Rev. D **66**, 125005 (2002) [hep-ph/0205034]; H. Abe, T. Higaki and T. Kobayashi, Prog. Theor. Phys. **109**, 809 (2003) [hep-th/0210025]; H. Abe, K. Choi and I. W. Kim, hep-th/0405100.

[6] Z. Chacko and M. A. Luty, JHEP **0105**, 067 (2001) [hep-ph/0008103].

[7] F. Borzumati and A. Masiero, Phys. Rev. Lett. **57**, 961 (1986);J. Hisano, T. Moroi, K. Tobe, M. Yamaguchi and T. Yanagida, Phys. Lett. B **357**, 579 (1995) [hep-ph/9501407]. J. Hisano, T. Moroi, K. Tobe and M. Yamaguchi, Phys. Rev. D **53**, 2442 (1996) [hep-ph/9510309];J. Hisano, D. Nomura and T. Yanagida, Phys. Lett. B **437**, 351 (1998) [hep-ph/9711348].

[8] J. Hisano and D. Nomura, Phys. Rev. D **59**, 116005 (1999) [hep-ph/9810479].

[9] J. A. Casas and A. Ibarra, Nucl. Phys. B **618**, 171 (2001) [hep-ph/0103065];S. Baek, T. Goto, Y. Okada and K. Okumura, Phys. Rev. D **64**, 095001 (2001)[hep-ph/0104146].

[10] A. Ceresole and G. Dall'Agata, Nucl. Phys. B **585**, 143 (2000) [hep-th/0004111]; R. Altendorfer, J. Bagger and D. Nemeschansky, Phys. Rev. D **63**, 125025 (2001) [hep-th/0003117]; T. Gherghetta and A. Pomarol, Nucl. Phys. B **586**, 141 (2000) [hep-ph/0003129]; A. Falkowski, Z. Lalak and S. Pokorski, Phys. Lett. B **491**, 172 (2000) [hep-th/0004093].

[11] K. Choi, H. D. Kim and Y. W. Kim, JHEP **0211**, 033 (2002) [hep-ph/0202257]; JHEP **0303**, 034 (2003) [hep-ph/0207013].

[12] L. J. Hall, Y. Nomura, T. Okui and S. J. Oliver, Nucl. Phys. B **677**, 87 (2004) [hep-th/0302192]; J. Bagger and D. Belyaev, JHEP **0306**, 013 (2003) [hep-th/0306063]; J. Bagger and M. Redi, Phys. Lett. B **582**, 117 (2004) [hep-th/0310086].

[13] N. Arkani-Hamed, T. Gregoire and J. Wacker, JHEP **0203**, 055 (2002) [hep-th/0101233]; D. Marti and A. Pomarol, Phys. Rev. D **64**, 105025 (2001) [hep-th/0106256]; A. Hebecker, Nucl. Phys. B **632**, 101 (2002) [hep-ph/0112230].

[14] G. v. Gersdorff, L. Pilo, M. Quiros, D. A. J. Rayner and A. Riotto, Phys. Lett. **B580**, 93 (2004) [hep-ph/0305218].

[15] Z. Chacko, M. A. Luty and E. Ponton, JHEP **0007**, 036 (2000) [hep-ph/9909248].

[16] T. Kobayashi and O. Vives, Phys. Lett. B **506**, 323 (2001) [hep-ph/0011200]; G. G. Ross and O. Vives, Phys. Rev. D **67**, 095013 (2003) [hep-ph/0211279]; S. F. King and I. N. R. Peddie, Nucl. Phys. B **678**, 339 (2004) [hep-ph/0307091]; K. S. Babu, Ts. Enkhbat and I. Gogoladze, Nucl. Phys. B **678**, 233 (2004) [hep-ph/0308093].

[17] A. B. Lahanas (Athens U.), D. V. Nanopoulos, Phys. Rept. **145**, 1 (1987); D. E. Kaplan, G. D. Kribs and M. Schmaltz, Phys. Rev. **D62**, 035010 (2000); Z. Chacko, M. A. Luty, A. E. Nelson and E. Ponton, JHEP **0001**, 003 (2000) [hep-ph/9911323]

[18] J. A. Casas, A. Lleyda and C. Munoz, Nucl. Phys. **B471**, 3 (1996) [hep-ph/9507294]; S. Komine, T. Moroi and M. Yamaguchi, Phys. Lett. B **507**, 224 (2001) [hep-ph/0103182].

[19] K. Choi, E. J. Chun and H. D. Kim, Phys. Lett. **B394**, 89 (1997) [hep-ph/9611293].

[20] L. J. Hall, V. A. Kostelecky and S. Raby, Nucl. Phys. B **267**, 415 (1986);R. Barbieri and L. J. Hall, Phys. Lett. B **338**, 212 (1994) [hep-ph/9408406];R. Barbieri, L. J. Hall and A. Strumia, Nucl. Phys. B **449**, 437 (1995) [hep-ph/9504373];J. Hisano, T. Moroi, K. Tobe and M. Yamaguchi, Phys. Lett. B **391**, 341 (1997) [Erratum-ibid. B **397**, 357 (1997)][hep-ph/9605296].

[21] R. Barbieri, L. J. Hall and A. Strumia, Nucl. Phys. B **445** (1995) 219 [hep-ph/9501334].

[22] F. Gabbiani and A. Masiero, Nucl. Phys. B **322**, 235 (1989); S. Baek, T. Goto, Y. Okada and K. Okumura, Phys. Rev. D **63**, 051701 (2001) [hep-ph/0002141]; *ibid.* **64**, 095001 (2001)[hep-ph/0104146]. T. Moroi, JHEP **0003**, 019 (2000) [hep-ph/0002208];Phys. Lett. B **493**, 366 (2000) [hep-ph/0007328];N. Akama, Y. Kiyo, S. Komine and T. Moroi, Phys. Rev. D **64**, 095012 (2001) [hep-ph/0104263];T. Goto, Y. Okada, Y. Shimizu, T. Shindou and M. Tanaka, Phys. Rev. D **66**, 035009 (2002) [hep-ph/0204081];hep-ph/0306093 .

[23] K. Inoue, A. Kakuto, H. Komatsu and S. Takeshita, Prog. Theor. Phys. **67**, 1889 (1982);*ibid.* **68** (1982) 927 [Erratum-ibid. **70** (1983) 330]; *ibid.* **71**, 413 (1984);H. Komatsu, INS-Rep.-469 (1983); K. Inoue, A. Kakuto and S. Takeshita, Prog. Theor. Phys. **71**, 348 (1984); A. Bouquet, J. Kaplan and C. A. Savoy, Phys. Lett. B **148**, 69 (1984);Nucl. Phys. B **262** (1985) 299.

[24] M. Dugan, B. Grinstein and L. J. Hall, Nucl. Phys. B **255**, 413 (1985);J. S. Hagelin, S. Kelley and T. Tanaka, Nucl. Phys. B **415**, 293 (1994);F. Gabbiani, E. Gabrielli, A. Masiero and L. Silvestrini, Nucl. Phys. B **477**, 321 (1996) [hep-ph/9604387];M. Misiak, S. Pokorski and J. Rosiek, Adv. Ser. Direct. High Energy Phys. **15**, 795 (1998) [hep-ph/9703442];M. Ciuchini *et al.*, JHEP **9810**, 008 (1998) [hep-ph/9808328].

[25] M. L. Brooks *et al.* [MEGA Collaboration], Phys. Rev. Lett. **83**, 1521 (1999) [hep-ex/9905013].

[26] L.M. Barkov *et al.*, "Search for the Decay $\mu^+ \to e^+\gamma$ Down to 10^{-14} Branching Ratio," Research Proporsal to Paul Scherrer Institut, 1999, available at http://meg.web.psi.ch/

[27] Letter of Intent for KEK Super B Factory, ed. by S. Hashimoto and M. Hazumi, available at http://belle.kek.jp/superb/

[28] K. Abe *et al.* [Belle Collaboration], Phys. Rev. Lett. **92**, 171802 (2004) [hep-ex/0310029].

[29] Y. Kuno and Y. Okada, Phys. Rev. Lett. **77** (1996) 434 [hep-ph/9604296] Y. Kuno, A. Maki and Y. Okada, Phys. Rev. D **55** (1997) 2517 [hep-ph/9609307] ; Y. Okada, K. Okumura and Y. Shimizu, Phys. Rev. D **61** (2000) 094001 [hep-ph/9906446] ; R. Kitano and Y. Okada, Phys. Rev. D **63** (2001) 113003 [hep-ph/0012040].

[30] K. Hagiwara *et al.* [Particle Data Group Collaboration], Phys. Rev. D **66**, 010001 (2002).

[31] A. J. Buras, S. Jager and J. Urban, Nucl. Phys. B **605** (2001) 600 [hep-ph/0102316].

[32] S. Bertolini, hep-ph/0206095; A. J. Buras and M. Jamin, JHEP **0401**, 048 (2004) [arXiv:hep-ph/0306217].

[33] C. Jessop, Report No. SLAC–PUB–9610.

[34] P. Gambino and M. Misiak, Nucl. Phys. B **611**, 338 (2001) [hep-ph/0104034].